Puzzlings in heavy quark energy loss

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Outline

- Soft radiation in QED and QCD
- Parton energy loss and
  Landau-Pomeranchuk-Migdal effect
- Energy loss of heavy quarks
- Beyond perturbative QCD: parton energy loss in sQGP
- Upper bound on the energy of the parton escaping from sQGP

Warm-up: particle production by a classical source

Consider the Lagrangian with a source $j(x)$ (antenna, jet, ...)

$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}m^2\phi^2 + j(x)\phi(x)$

The corresponding Euler-Lagrange equation is the Klein–Gordon equation:

$(\partial^\mu \partial_\mu + m^2)\phi = j(x)$ or $\left(\frac{\partial^2}{\partial t^2} - \nabla^2 + m^2\right)\phi = j(x)$.

Let us assume $j(x)$ is limited in time; then before the source is turned on, the field is free:

$\phi(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2\omega_p}} \left(a_p e^{ipx} + a_p^\dagger e^{-ipx}\right)$

After the action of the source, the solution will be given by

$$\phi(x) = \phi_0(x) +$$
$$i \int d^4y \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2\omega_p}} \theta(x^0 - y^0) \left(e^{-ip(x-y)} - e^{ip(x-y)}\right) j(y)$$

Introducing the Fourier transform $\tilde{j}(p) = \int d^4y e^{ipy} j(y)$, at $p^2 = m^2$, we see that the free field operators are changed:

$$a_p \to a_p + \frac{i}{\sqrt{2\omega_p}} \tilde{j}(p), \quad a_p^\dagger \to a_p^\dagger - \frac{i}{\sqrt{2\omega_p}} \tilde{j}^*(p)$$

cf superfluidity

The Hamiltonian then is

$$H = \int \frac{d^3p}{(2\pi)^3} \omega_p \left(a_p + \frac{i}{\sqrt{2\omega_p}} \tilde{j}(p) \right) \left(a_p^\dagger - \frac{i}{\sqrt{2\omega_p}} \tilde{j}^*(p) \right)$$

The energy of the vacuum after the source has been turned off is

$$\langle 0|H|0\rangle = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2} |\tilde{j}(p)|^2$$

The particle number is

$$\int dN = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2\omega_p} |\tilde{j}(p)|^2$$

Radiated field is coherent,
Poisson distribution in the number of quanta

Let us now consider instead of the scalar field with a source $j(x)$

$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}m^2\phi^2 + j(x)\phi(x)$

electromagnetic field with a source J_μ:

$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - J_\mu A^\mu$

The only difference from the scalar case is that photons are vector particles with polarization $e_\mu(\vec{p})$, and so we should replace

$$\phi \to A^\mu$$

and thus

$$\int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2\omega_p}} e^{ipx} \to \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2\omega_p}} e_\mu(\vec{p}) e^{ipx}$$

This leads to $J \to e^\mu J_\mu$.

In complete analogy with the previous case, we get for the total energy of produced photons (intensity of radiation J)

$$dJ = \frac{d^3p}{(2\pi)^3}\frac{1}{2}|e^\mu J_\mu(p)|^2$$

and for the total average number of photons emitted with frequencies $\omega_1 \leq \omega \leq \omega_2$:

$$\bar{N} = \int_{\omega_1 \leq \omega \leq \omega_2} \frac{d^3p}{(2\pi)^3}\frac{1}{2\omega_p}|e^\mu J_\mu(p)|^2$$

Here the charged classical source $J_\mu(p)$ in electrodynamics is called a **transition current**. It is conserved (charge conservation):

$$p^\mu J_\mu = 0$$

For a charge (e.g. an electron of mass m) undergoing a scattering from the initial 4-momentum k_1 to the final 4-momentum k_2 the transition current is

$$J_\mu = e\sqrt{4\pi}\left(\frac{k_{1\mu}}{pk_1} - \frac{k_{2\mu}}{pk_2}\right);$$

clearly, $p^\mu J_\mu = 0$.

Non-relativistic limit: $|\vec{k}_1|, |\vec{k}_2| \ll m$, $pk_1 \simeq \omega m$, $pk_2 \simeq \omega m$; then

$$\vec{J} = e\sqrt{4\pi}\,\frac{1}{\omega}\,(\vec{v}_1 - \vec{v}_2), \quad J_0 \to 0,$$

where $\vec{v}_i = \vec{p}_i/m$.

If $\vec{v}_1 = \vec{v}_2$ radiation vanishes (**no scattering, no radiation**)

Since photons are vector particles, we have a polarization vector e^μ, and the probability of radiation will depend on $(e^\mu J_\mu)$;

The probability of photon radiation is proportional to the differential cross section of charge scattering, and since

$$\bar{N} = \int_{\omega_1 \le \omega \le \omega_2} \frac{d^3p}{(2\pi)^3} \frac{1}{2\omega_p} |e^\mu J_\mu(p)|^2,$$

the differential cross section of photon emission is

$$d\sigma_{brems} \sim d\sigma_s \; e^2 \; |\vec{e}(\vec{v}_1 - \vec{v}_2)|^2 \; d\Omega \; \frac{d\omega}{\omega}$$

Note that

$$\frac{d\sigma_{brems}}{d\omega} \to \infty \text{ when } \omega \to 0$$

– this is the **infrared catastrophe.**

The total cross section for photon production is

$$d\sigma_t \sim d\sigma_s \; \alpha \ln\left(\frac{\omega_{max}}{\omega_{min}}\right)$$

When $\alpha \ln\left(\frac{\omega_{max}}{\omega_{min}}\right) \simeq 1$, we deal with the radiation of a very large number of photons = classical electromagnetic field.

$\omega_{max} = k_0 \equiv E$; this is how **logarithms of energy** $\alpha \ln E$ appear.

High energies = classical radiation; $\bar{N} \sim \frac{1}{\alpha}$.

CGC

Soft radiation off a high energy particle

(electron, quark, quasi-particle, ...)

Consider now the case of a high–energy particle, when $k_{10} \gg m$, and soft radiation, $p_0/k_{10} \ll 1$.

For a photon emitted at an angle θ_1, we have

$$(pk_1) = p_0 k_{10} - p_0 |\vec{k}_1| \cos\theta_1$$

Using $k_{10} = \sqrt{m^2 + |\vec{k}_1|^2} \simeq |\vec{k}_1| + \frac{m^2}{2|\vec{k}_1|}$, we get

$$(pk_1) \simeq p_0 |\vec{k}_1| \left(1 - \cos\theta_1 + \frac{m^2}{2|\vec{k}_1|^2} \right)$$

For small angles $1 - \cos\theta \simeq \theta^2/2$, and introducing $\theta_0^2 = m^2/|\vec{k}_1|^2$

we get

$$(pk_1) \simeq \frac{|\vec{k}_1|p_0}{2}\left(\theta_1^2 + \theta_0^2\right)$$

When θ_1, θ_0 are small, J is large -

photons are emitted mainly at small angles relative to the direction of the electron.

The numerator $(ek_1) = -|\vec{k}_1|\sin\theta_1 \simeq -|\vec{k}_1|\theta_1$ (the photon polarization lies in the scattering plane $\{\vec{k}_1, \vec{k}_2\}$

The probability to radiate is thus

$$\sim \frac{e^2}{p^2}\left(\frac{\theta_1}{\theta_1^2 + \theta_0^2} - \frac{\theta_2}{\theta_2^2 + \theta_0^2}\right)^2$$

Note that the large electron momentum has canceled out - no dependence, except through $\theta_0 = m^2/2|\vec{k}_1|^2$ –

narrowing at high energy.

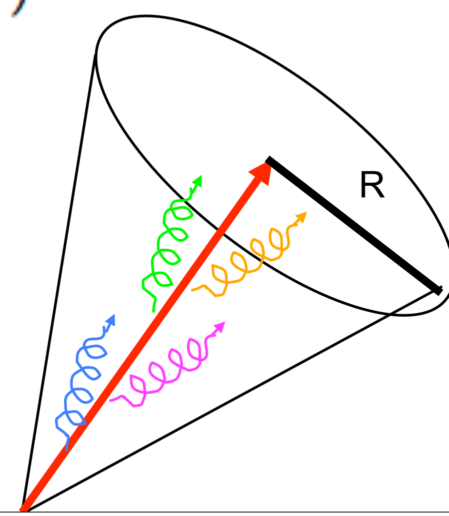
If scattering angle θ_s is small, $\theta_s \ll \theta_1, \theta_2$ then we have almost complete cancelation - no scattering, no radiation.

If $\theta_s \ll \theta_0$, also no radiation. But if $\theta_s \gg \theta_0$, radiation is strong in two cases: $\theta_1 \simeq \theta_s \gg \theta_2$ or viceversa –

two identical narrow cones of radiation directed along $\vec{k_1}$ and $\vec{k_2}$

Consider the photon distribution in one of these cones:

$$d\sigma_{brems} = d\sigma_s \; \frac{4e^2}{p^2} \; \frac{\theta_1^2}{(\theta_1^2 + \theta_0^2)^2} \; \frac{d^4p \; \delta(p^2)}{(2\pi)^3}$$



For small angle emission, $\theta_1, \theta_0 \ll 1$ the differential cross section is given by

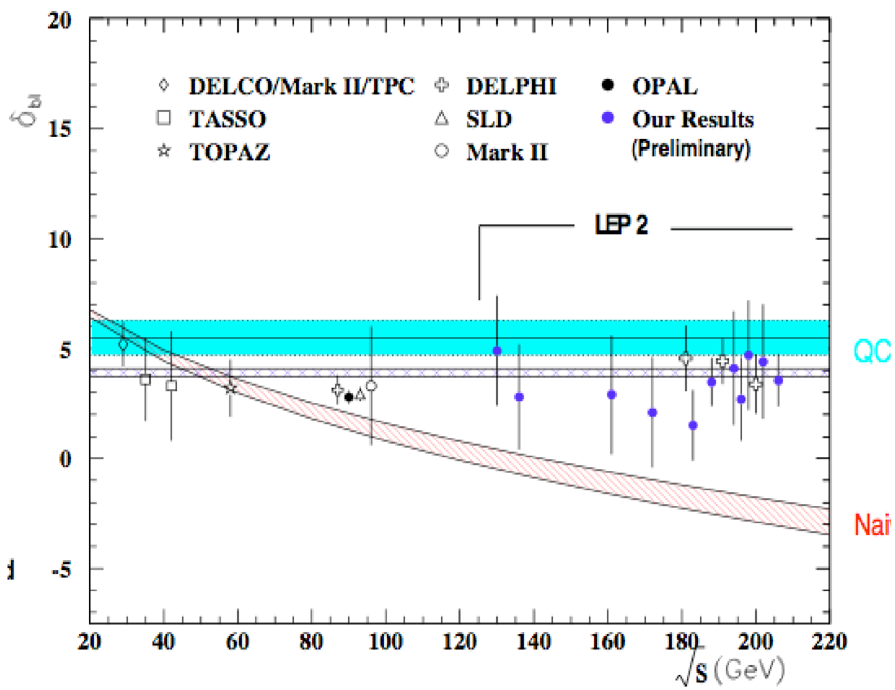
$$d\sigma_{brems} = d\sigma_s \, \frac{\alpha}{\pi} \frac{dp_0}{p_0} \, \frac{\theta_1^2 d\theta_1^2}{(\theta_1^2 + \theta_0^2)^2}.$$

Note that for massive particles with $\theta_0 = m^2/2|\vec{k}_1|^2 \neq 0$, the radiation in the forward cone vanishes – **dead cone effect**:

massive particles radiate less

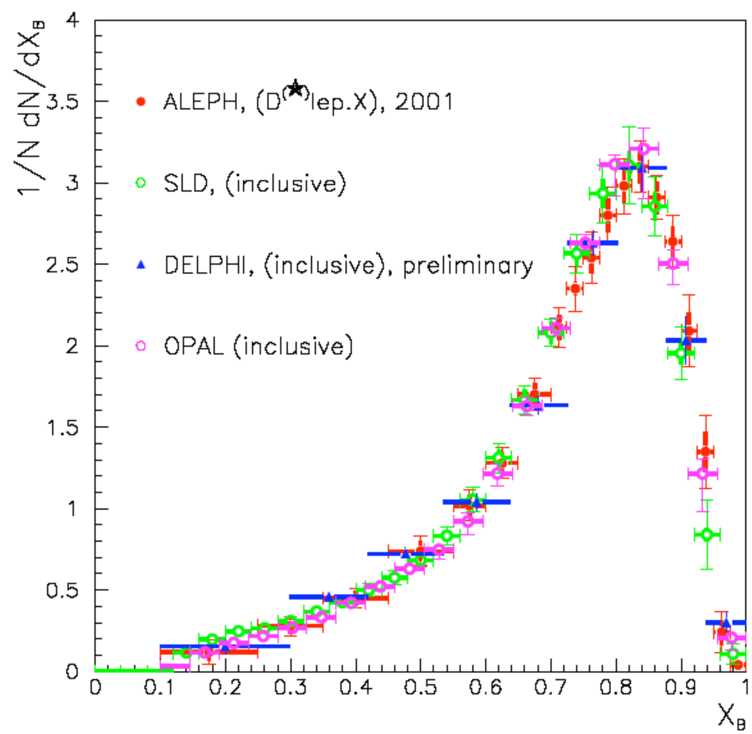
Can be shown to be a consequence of causality
for a source propagating with v<c

Heavy quarks fragment differently



OPAL Collaboration

Heavy quarks produce a larger number of particles



and carry a larger fraction of jet momentum

Assuming large–angle scattering, $\theta_s \sim 1$ and Integrating over the photon energies and photon angles in the cone $\theta_0 < \theta_1 < \theta_s$ we get

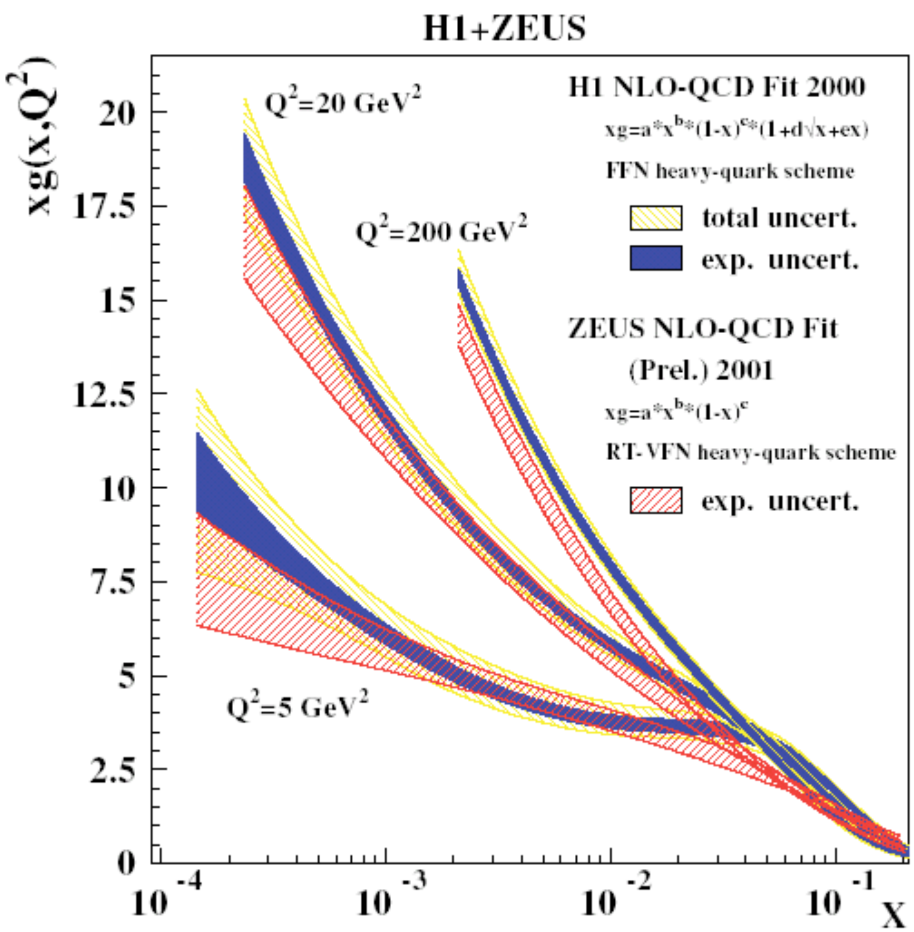
$$d\sigma_{brems} = d\sigma_s \, \frac{\alpha}{\pi} \ln \frac{p_{0max}}{p_{0min}} \, \ln \frac{\theta_s^2}{\theta_0^2} \simeq$$

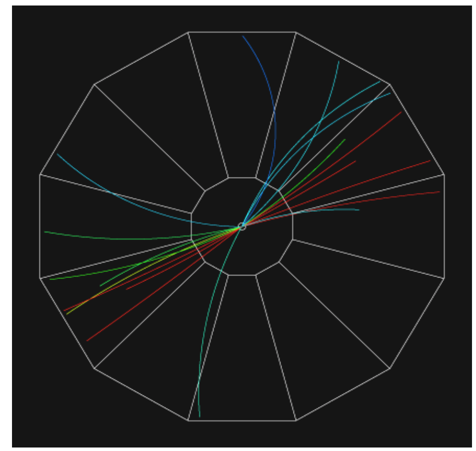
$$= d\sigma_s \, \frac{\alpha}{\pi} \ln \frac{|\vec{k}_1|}{p_{0min}} \, \ln \frac{\theta_s^2 |\vec{k}_1|^2}{m^2}$$

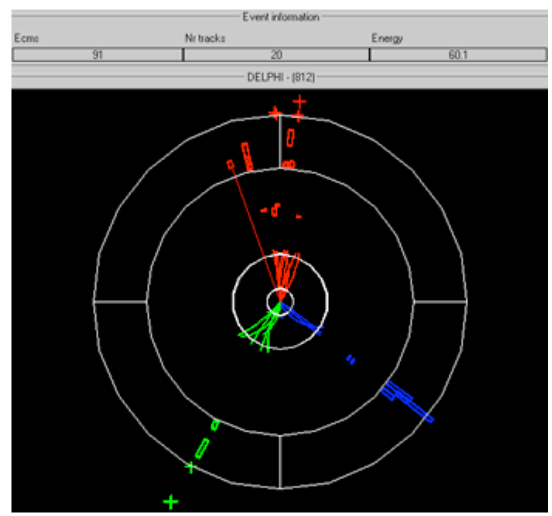
"perturbative logs" – longitudinal and transverse.

Their "resummation" is a major tool in QCD and QED perturbation theory at high energies; it gives rise e.g. to the evolution of parton densities

Gluon emission as seen in the structure functions and in the jet structure







Qualitative picture of gluon emission

Formation time: the emission of a gluon takes a time which can be estimated as follows:

$$t_{\text{fluct}} \sim \frac{E_1}{|m^2 - (p_1 - k)^2|} = \frac{E_1}{2p_1 k} \sim \frac{1}{\omega\Theta^2} \simeq \frac{\omega}{k_\perp^2}\,.$$

$$k_\perp \approx \omega\Theta \ll k_{||} \approx \omega.$$

Quanta with small transverse momenta, large energy take a long time to form!

The photon (gluon) field is localized within

$$\lambda_{||} \sim \omega^{-1}, \; \lambda_{\perp} \sim k_{\perp}{}^{-1}$$

Now, suppose the charge got scattered at an angle θ_s: during the formation time, the charge is localized within

$$\Delta r_{||} \sim \left| v_{2||} - v_{1||} \right| \cdot t_{\text{fluct}} \sim \Theta_s^2 \cdot \frac{1}{\omega\Theta^2} = \left(\frac{\Theta_s}{\Theta}\right)^2 \lambda_{||} \;\; \Leftrightarrow \;\; \lambda_{||} \; ;$$

$$\Delta r_{\perp} \sim \quad c\Theta_s \cdot t_{\text{fluct}} \sim \Theta_s \cdot \frac{1}{\omega\Theta^2} \quad = \left(\frac{\Theta_s}{\Theta}\right) \lambda_{\perp} \;\; \Leftrightarrow \;\; \lambda_{\perp} \, .$$

Radiation occurs if the scattered charge “gets away” from the photon field during the formation time; at large scattering angles, it happens for all emission angles - 2 cones. At small scattering angle, large angle emission is suppressed: **angle ordering** in QCD cascades

Gluon radiation in QCD

All differences stem from the non-Abelian nature of QCD:

$$\mathcal{L} = \bar{\psi}(i\hat{D} - m)\psi - \frac{1}{4}\left(F^a_{\mu\nu}\right)^2$$

where

$$\hat{D} = \gamma^\mu D_\mu; \quad D_\mu = \partial_\mu - igA^a_\mu t^a,$$

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + gf^{abc}A^b_\mu A^c_\nu$$

Euler-Lagrange equation of motion (non-Abelian Maxwell equation):

$$\partial^\mu F^a_{\mu\nu} + gf^{abc}A^{b\mu}F^c_{\mu\nu} = -gj^a_\nu$$

or, in the compact form

$$D^{\mu} F^{a}_{\mu\nu} = -g j^{a}_{\nu}$$

The color current of quarks is

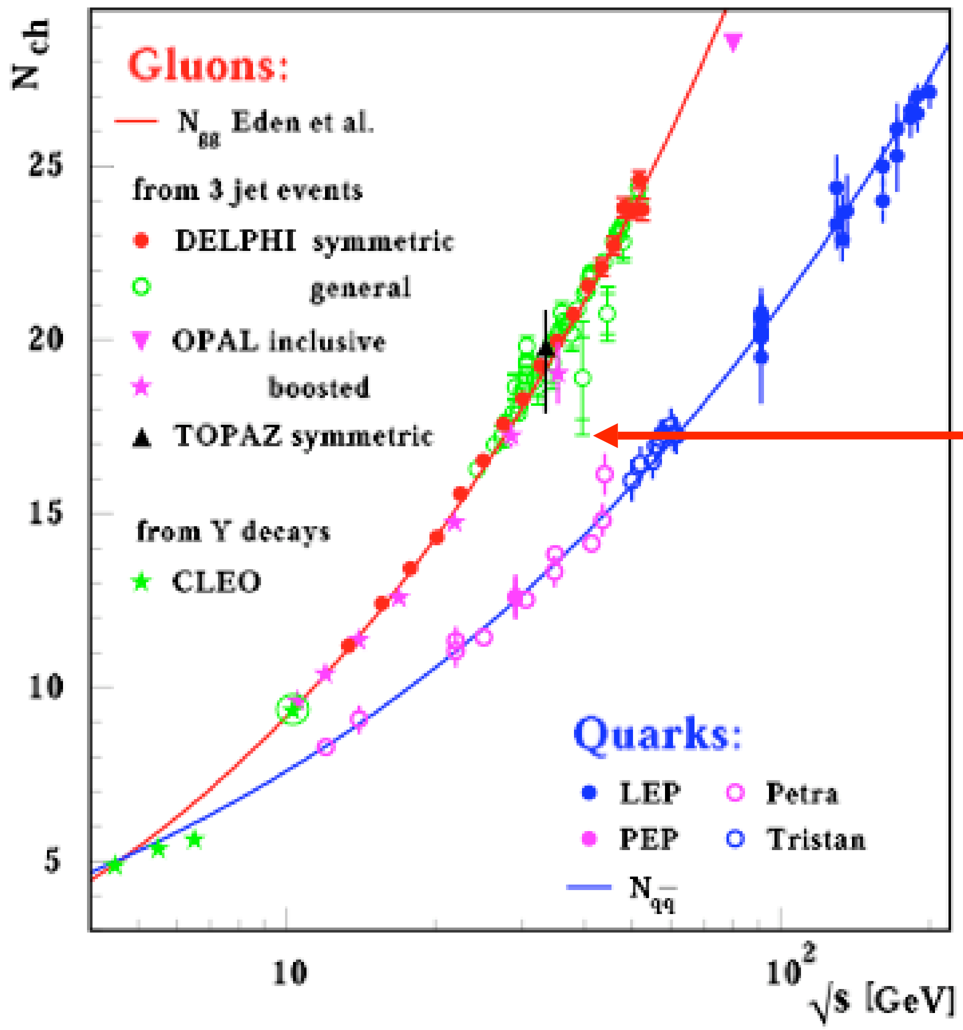
$$j^{a}_{\nu} = \bar{\psi} \gamma_{\nu} t^{a} \psi$$

The generators of color rotations obey commutation relation

$$[t^{a}, t^{b}] = i f^{abc} t^{c}$$

where f^{abc} are the **structure constants** of the $SU(N)$ group.

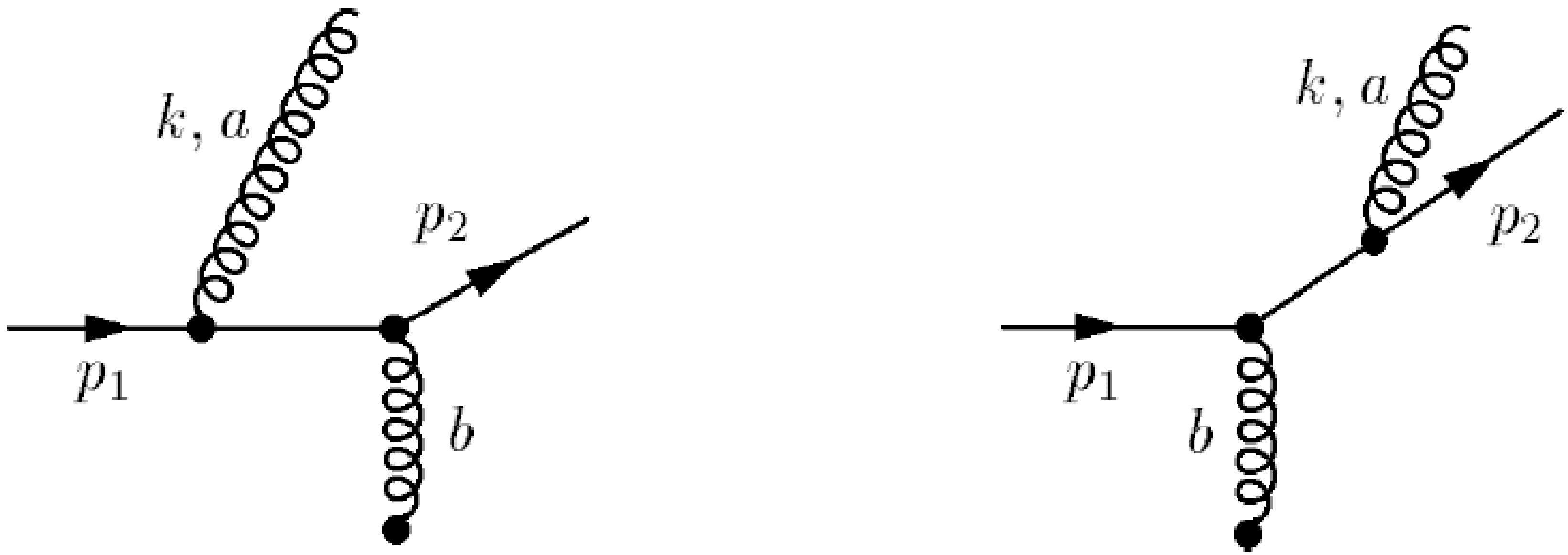
Quark and gluon jets fragment differently: color charge matters



The difference in hadron multi-plicities becomes visible at large momenta

Tagging gluon jets by g->ccbar and quark jets by leading charmed hadrons?

Soft gluon radiation in QCD



$$j^{\mu} = \left[t^b\, t^a \left(\frac{p_1^{\mu}}{(p_1 k)} \right) - t^a\, t^b \left(\frac{p_2^{\mu}}{(p_2 k)} \right) \right]$$

Introducing the abbreviation $A_i = \frac{p_i^\mu}{(p_i k)}$, we apply the standard decomposition of the product of two triplet color generators,

$$t^a t^b = \frac{1}{2N_c}\delta_{ab} + \tfrac{1}{2}\left(d_{abc} + \mathrm{i} f_{abc}\right)\, t^c \,,$$

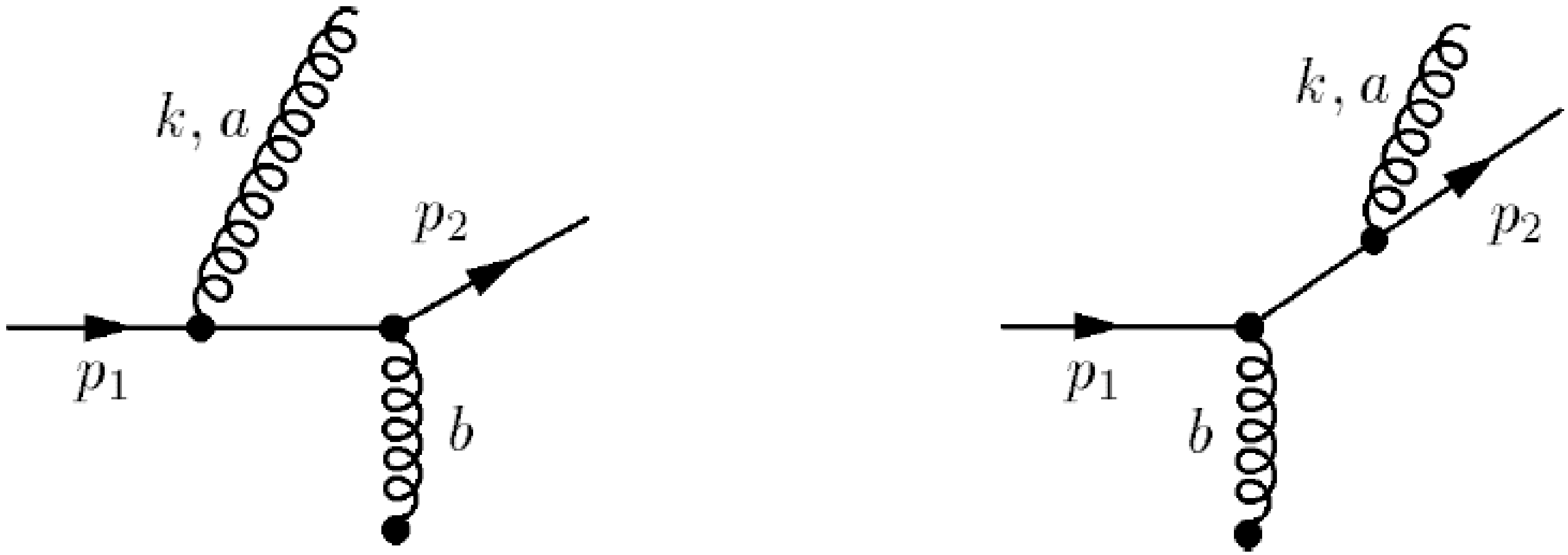
we get

$$\mathrm{d}N \;\propto\; \frac{1}{C_F}\sum_{\text{color}} j^\mu \cdot (j_\mu)^* = \left(\frac{1}{2N_c} + \frac{N_c^2-4}{4N_c}\right)(A_1 - A_2)\cdot(A_1 - A_2)$$
$$+\; \frac{N_c}{4}\,(A_1 + A_2)\cdot(A_1 + A_2)\,.$$

which simplifies to

$$\mathrm{d}N \;\propto\; C_F\,(A_1 - A_2)\cdot(A_1 - A_2) \;+\; N_c\, A_1 \cdot A_2$$

Yu.Dokshitzer



$$dN \propto C_F\,(A_1 - A_2)\cdot(A_1 - A_2) + N_c\,A_1\cdot A_2$$

Qualitative picture: 2 narrow bremsstrahlung cones along the directions of scattered quarks; one broad cone around the direction of the exchanged gluon - dominates the region of large emission angles

Multiple scattering in the medium: Landau-Pomeranchuk-Migdal effect

The gluon during its formation time

$$t_{form} \simeq \frac{\omega}{k_\perp^2}$$

accumulates a typical transverse momentum

$$k_\perp^2 \simeq \mu^2 \; \frac{t_{form}}{\lambda},$$

where λ is the mean free path and μ^2 is the characteristic momentum squared acquired in a single scattering.

This is a random walk with the average number of coherent scatterings

$$N_{coh} = \frac{t_{form}}{\lambda} = \sqrt{\frac{\omega}{\mu^2 \lambda}}$$

For sufficiently high energies the coherent length exceeds the mean free path, $N_{coh} > 1$, and the standard independent Bethe-Heitler emission gets suppressed:

$$A_1 - A_2 + A_2 - A_3 + A_3 + \dots - A_N$$

This leads to the suppression in the emission spectrum:

$$\frac{dW}{d\omega dz} = \left(\frac{dW}{d\omega dz}\right)^{\mathrm{BH}} \cdot \frac{1}{N_{coh}} = \frac{\alpha_s C_R}{\pi\omega\,\lambda} \cdot \sqrt{\frac{\mu^2\,\lambda}{\omega}} = \frac{\alpha_s C_R}{\pi\omega}\sqrt{\frac{\hat{q}}{\omega}}.$$

where we have introduced the **transport coefficient**

$$\hat{q} \equiv \rho \int \frac{d\sigma}{dq^2}\, q^2\, dq^2,$$

BDMPS; Zakharov; GLV; GWW; AMY; ...

The transverse momentum and the energy of the emitted gluon are related by

$$k_\perp^2 \simeq \sqrt{\hat{q}\,\omega}.$$

Therefore the angular distribution of the emitted gluon is concentrated in a characteristic energy- and medium- dependent emission angle

$$\theta \simeq \frac{k_\perp}{\omega} \sim \left(\frac{\hat{q}}{\omega^3}\right)^{1/4}$$

Heavy quarks: as we already discussed, the emission at small angles is suppressed:

$$dP = \frac{\alpha_s\, C_F}{\pi}\,\frac{d\omega}{\omega}\,\frac{k_\perp^2\, dk_\perp^2}{(k_\perp^2 + \omega^2\theta_0^2)^2}, \qquad \theta_0 \equiv \frac{M}{E},$$

This suppression results in the modification of the gluon radiation spectrum:

$$I(\omega) = \omega \frac{dW}{d\omega} = \frac{\alpha_s\, C_F}{\pi} \sqrt{\frac{\omega_1}{\omega}}\, \frac{1}{(1 + (\ell\, \omega)^{3/2})^2},$$

where

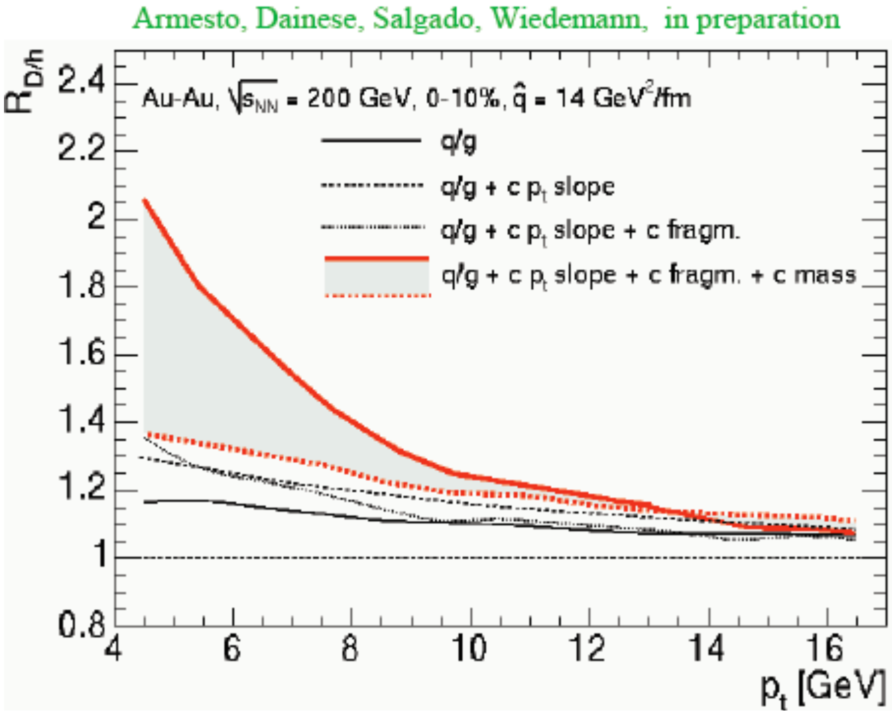
$$\ell \equiv \hat{q}^{-1/3} \, \left(\frac{M}{E}\right)^{4/3} .$$

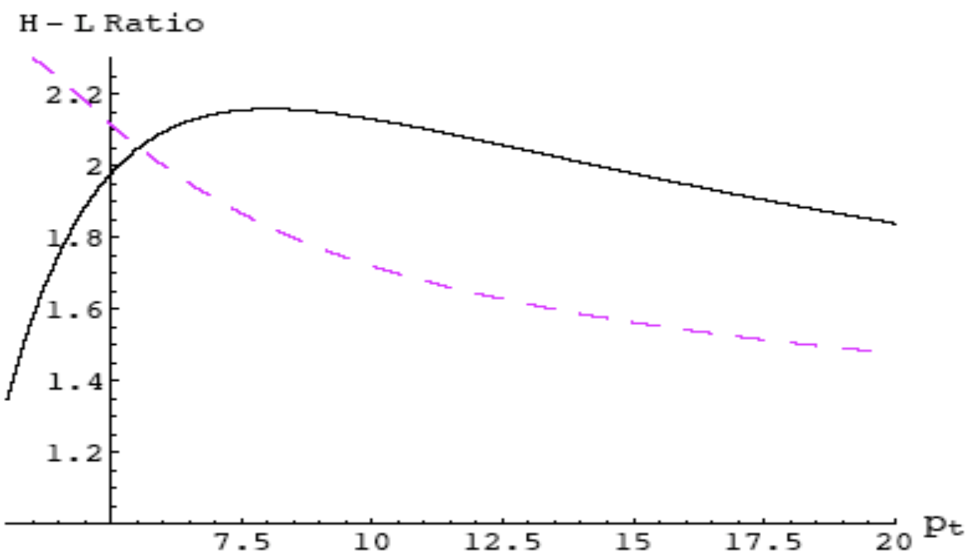
Yu.L.Dokshitzer and DK, ‘01

“Dead cone” effect in the induced radiation:
Energy loss of heavy quarks should be suppressed relative to the energy loss of light quarks

Perturbative energy loss hierarchy:

E (heavy quark) < E (light quark) < E (gluon)





Enhancement of
the D/h ratio as
a signature of the radiative
energy loss in the QGP

Heavy quark colorimetry of QCD matter

col·or·im·e·try *noun*
col·or·im·e·ter *noun* **:**
an instrument or device for determining and specifying colors; *specifically* **:** one used for chemical analysis by comparison of a liquid's color with standard colors

Merriam-Webster Dictionary

The propagation of heavy quarks in QCD matter is strongly affected by the interplay of the “dead cone” and quantum interference effects (LPM) at energies up to

$$E \leq M\sqrt{\hat{q}L^3}$$

(a consequence of quantum mechanics & causality)

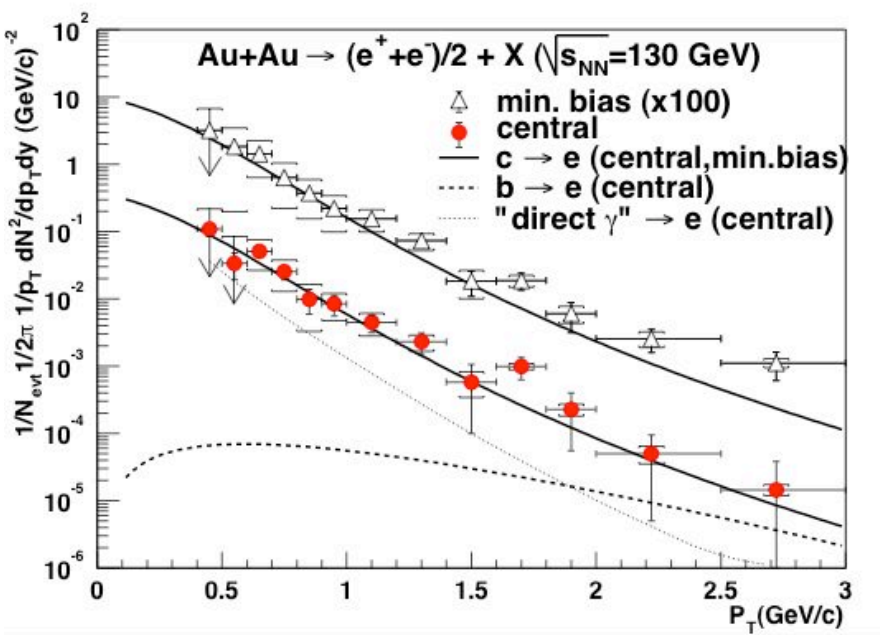
For heavy quarks the induced gluon radiation should be suppressed; is it?

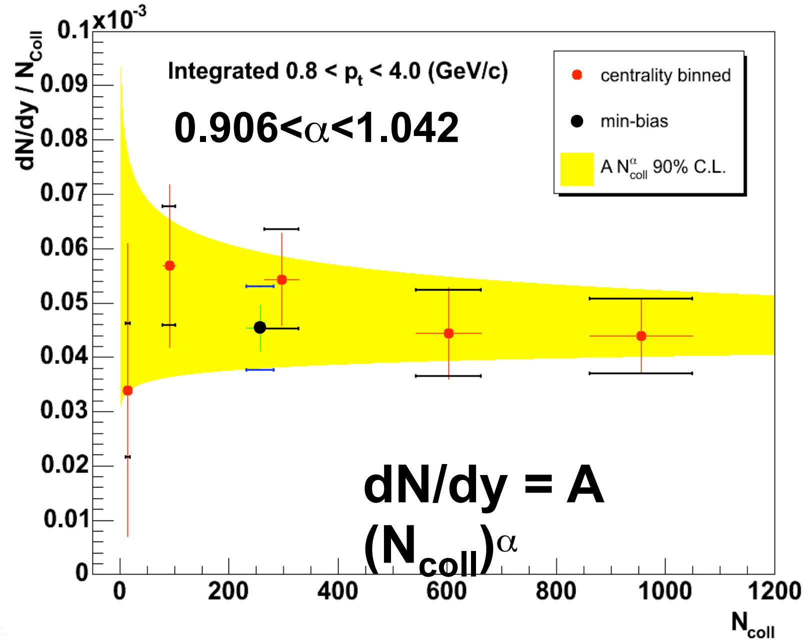
Recent work:

M.Djordjevic, M.Gyulassy '03-

B.Zhang, E.Wang, X.-N. Wang'04

N.Armesto, C.Salgado, U.Wiedemann'04-

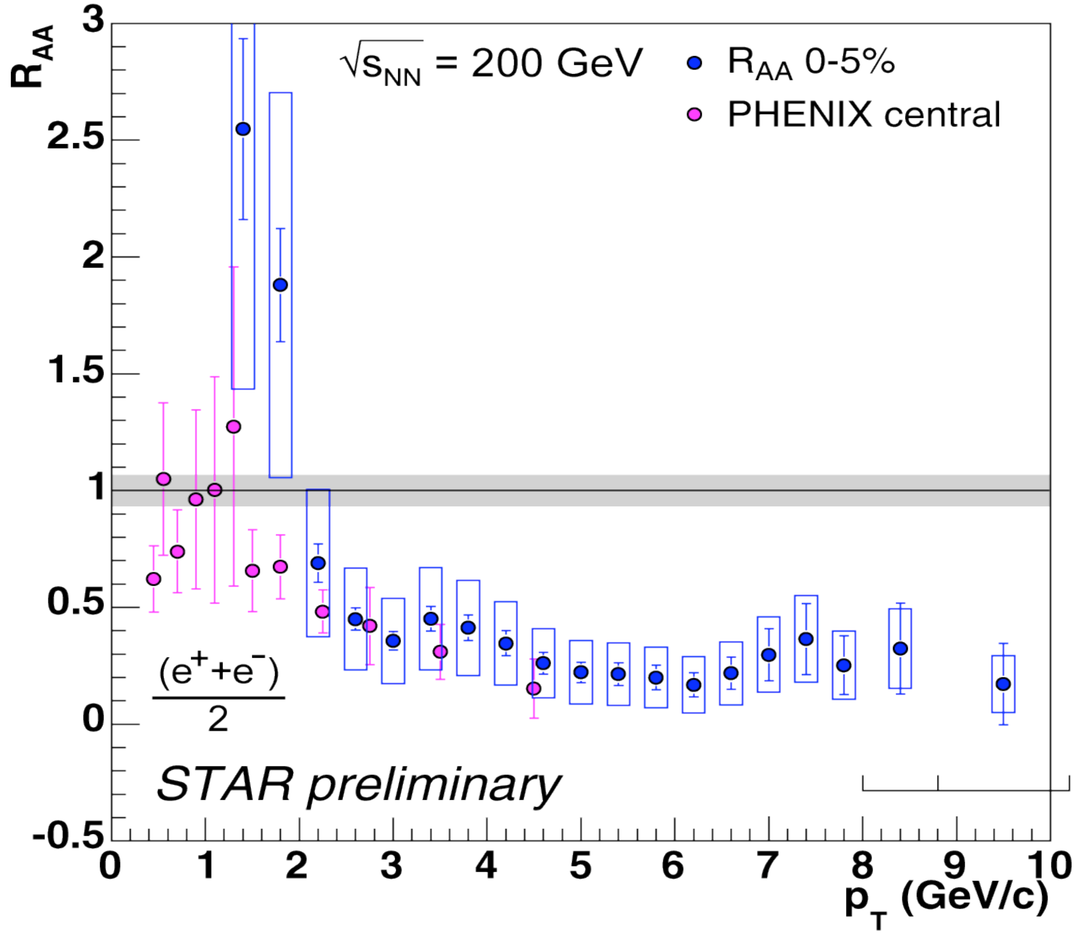




Data from PHENIX

AuAu collisions: charm is quenched!?

a serious problem for the naïve radiative energy loss scenario?



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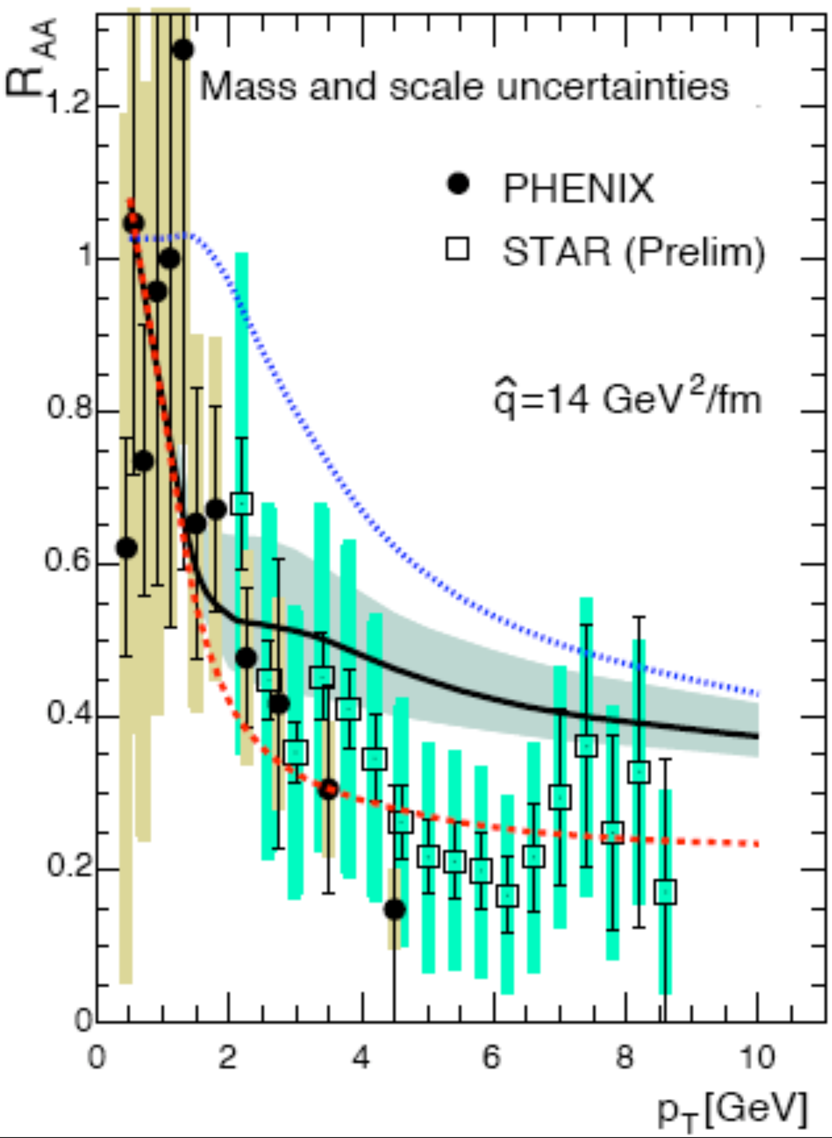
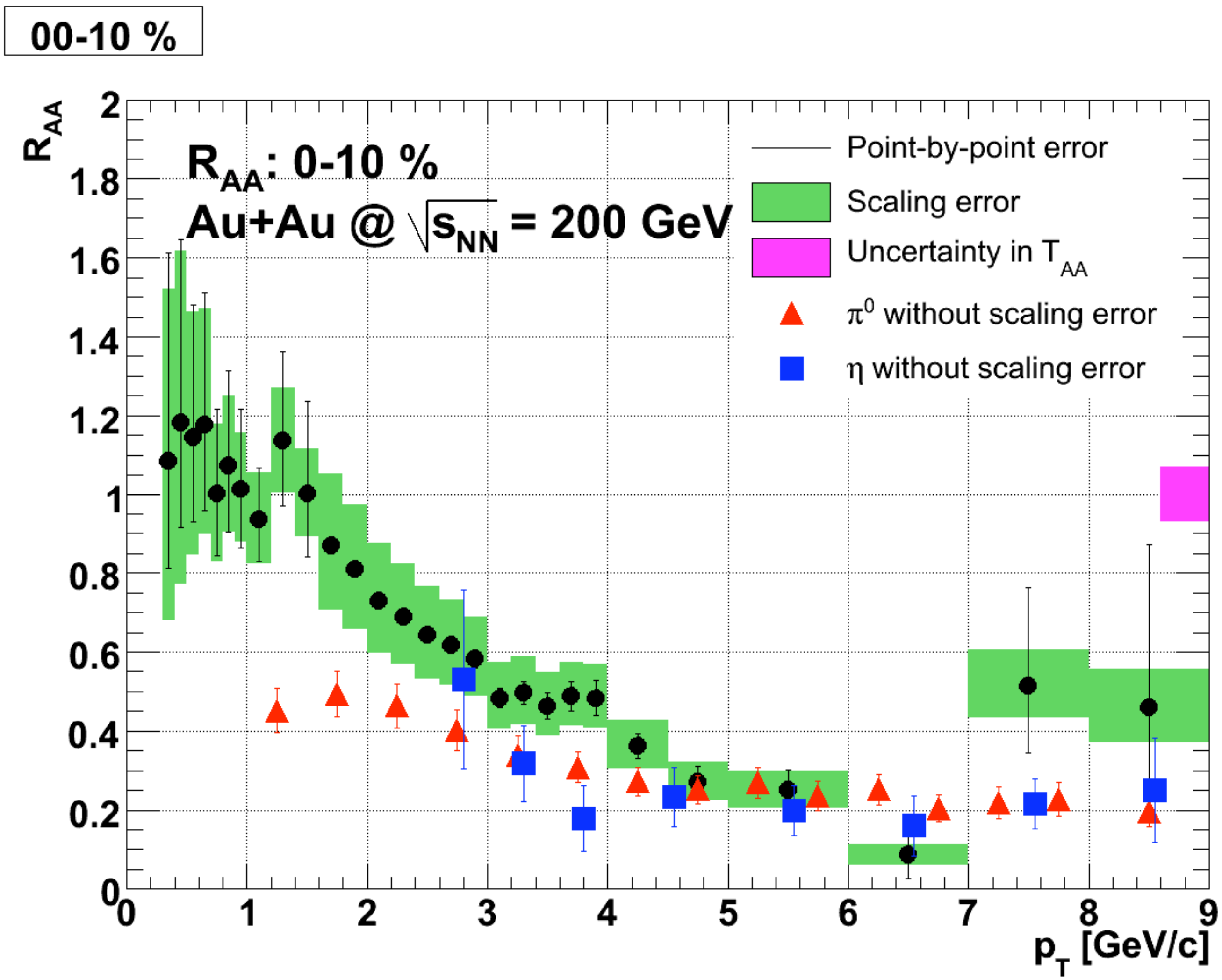


N.Armesto. M.Cacciari, A.Dainese, C.Salgado,U.Wiedemann, hep-ph/0511257

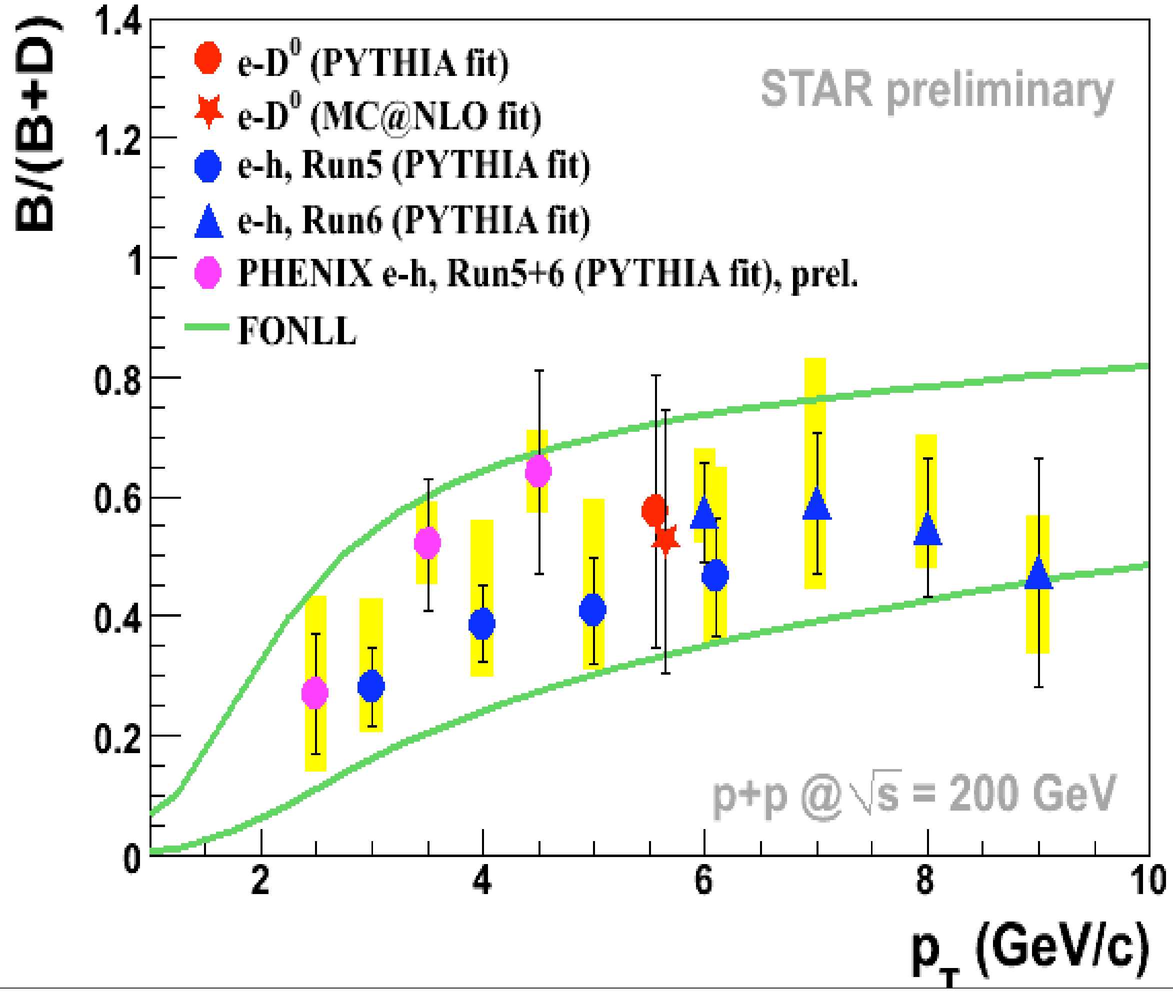
Need to separate b and c contributions!

Heavy quarks are suppressed:



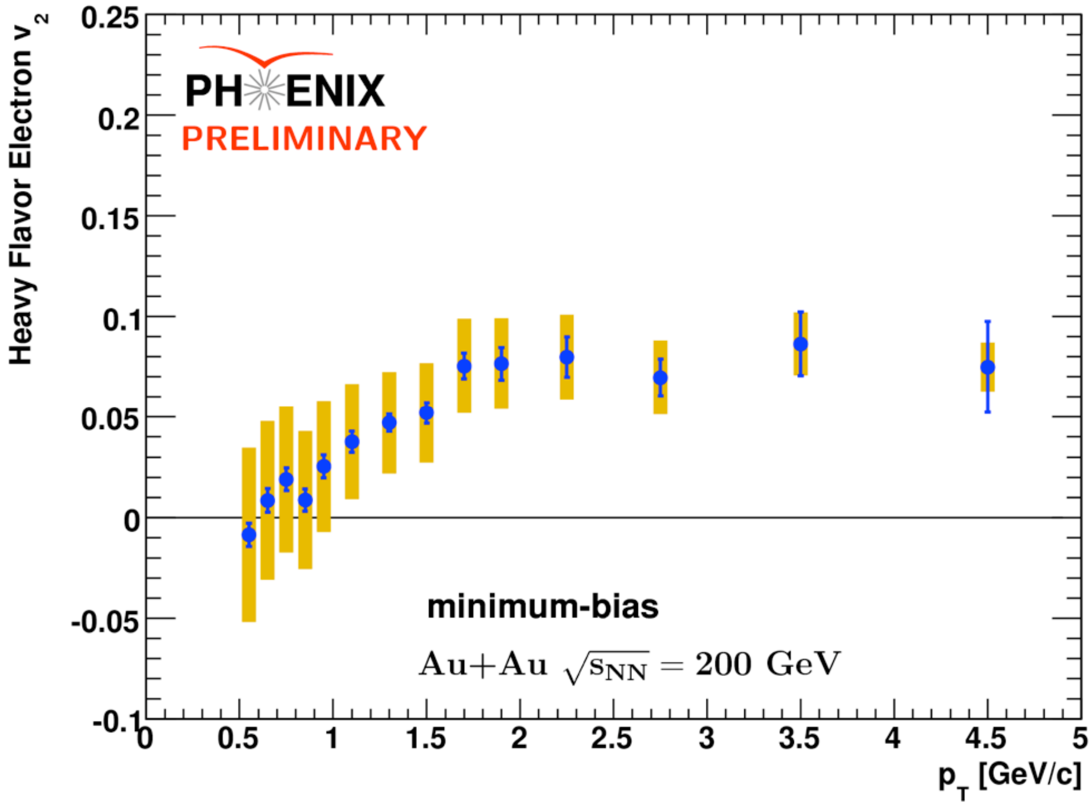
QM'08

Even b-quarks seem to be suppressed:



QM '08

Moreover, c and b-quarks even seem to flow:



QM '08

Is there a way to understand this?

Beyond perturbative energy loss?

In the perturbative approach, the energy lost is assumed to be small compared to the total energy of the parton.

The energy loss due to radiation in BDMPS approach is given by

$$E_{rad} = \frac{\alpha_s N_c}{4} \ \hat{q} L^2;$$

For $\hat{q} = 1$ GeV^2/fm, $L = 5$ fm, $\alpha_s = 0.3$, we get

$$E_{rad} \simeq 30 \text{ GeV};$$

at RHIC range of parton energies, this is by no means small !

Likewise, the rescattering of the radiated gluon is treated in the eikonal, small-angle approximation.

This requires the angle be small:

$$\theta \simeq \frac{p_\perp}{xp} \simeq \frac{\sqrt{\hat{q}L}}{xp} \ll 1$$

For $\hat{q} = 1\ \mathrm{GeV}^2/\mathrm{fm}$, $L = 5$ fm, $\sqrt{\hat{q}L} \simeq 2$ GeV.

Assuming $x \sim 0.1$, $\theta \sim 0.1$ requires the momentum of the jet be $p \sim 200$ GeV.

Can we go beyond the eikonal, perturbative treatments?

Classical Electrodynamics:

yes, the answer is given by the Liénard formula:

$$E_{rad}(T) = \frac{2e^2}{3} \int_{-\infty}^{T} dt \ \frac{\vec{a}^2 - (\vec{v} \times \vec{a})^2}{(1 - \vec{v}^2)^3}.$$

Energy loss from different segments of the path is simply added: the consequence of the fact that Maxwell equations are linear –

The field radiated at any time is independent of the previously radiated field, and is simply added to it.

Yang-Mills equations are highly non-linear, and the radiated field in general is a non-local functional of the trajectory.

However in AdS/CFT in the large N limit the answer for the energy radiated by a heavy quark is given by [Mikhailov '03, ...]:

$$E_{rad}(T) = \frac{\sqrt{\lambda}}{2\pi} \int_{-\infty}^{T} dt \; \frac{\vec{a}^2 - (\vec{v} \times \vec{a})^2}{(1 - \vec{v}^2)^3},$$

which differs from electrodynamics only by the substitution $2e^2/3 \leftrightarrow \sqrt{\lambda}/2\pi$.

The linear form of this formula is now not due to the linear equations of motion, but due to the integrability at large N.

If we <u>assume</u> that the energy loss indeed is given by the Liénard formula, we arrive at quite remarkable consequences...

The Liénard energy loss in external fields grows as a square of the energy:

$$E_{rad} \sim \frac{1}{1 - v^2} = \frac{E^2}{m^2}$$

(in AdS/CFT, this is a consequence of graviton emission; in classical electrodynamics - tensor $F_{\mu\nu}$).

Therefore, energy loss satisfies the equation

$$-\frac{dE}{dx} = k(x) E^2$$

Integrate:

$$\frac{1}{E_f} = \frac{1}{E_0} + \int_{-\infty}^{x} dx \; k(x)$$

As $E_0 \to \infty$, the final energy of the parton E_f tends to a constant value independent of E_0.

[I. Ya. Pomeranchuk, 1939]

To estimate how big this limit is, let us first parameterize the strength of fields. Classical treatment fails when the field reaches the value

$$F_c \equiv \frac{m^2}{g^3}$$

Let us measure the field in sQGP in units of the critical field computed for charm quark:

$$F = z\ F_c; \quad 0 < z < 1.$$

we expect that z is in fact quite close to 1.

The upper bound on the energy of the parton escaping from sQGP is given by

$$\frac{1}{E_{crit}} = \frac{2}{3m^2}\ \frac{g^2}{m^2}\ g^2 F^2\ L.$$

We thus get a universal (independent of initial energy) upper bound on the energy of charm quark [DK, to appear]:

$$E_{crit} = \frac{3}{2}\ \frac{g^2}{z^2 L},$$

or

$$E_{crit} = \frac{9}{8\pi}\ \frac{\sqrt{\lambda}}{z^2 L},$$

For a quark of flavor f, this becomes

$$E^f_{crit} = \frac{3}{2}\ \left(\frac{m_f}{m_c}\right)^2\ \frac{g^2}{z^2 L},$$

Estimate: $L = 5$ fm, $\alpha_s = 0.5$, $z \simeq 1$:

$E^c_{crit} \simeq 0.35$ GeV; $E^b_{crit} \simeq 3.2$ GeV.

This means that all charm quarks and a large fraction of bottom quarks will be stopped in the medium – **no** c or b quarks with energy above E_{crit} except the fraction coming from the surface of the fireball.

Very little room for observing a medium-modified jet – either there is a unmodified jet from the surface, or no jet at all

Summary

- The RHIC results on the production of charm and beauty are hard (impossible?) to reconcile with pQCD approach to the energy loss
- If AdS/CFT provides a qualitatively correct guidance to what happens at strong coupling and strong fields, there should exist a universal upper bound on the energy of partons escaping from sQGP