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Dynamics of the Landau–Pomeranchuk–Migdal effect in Au + Au collisions at $\sqrt{s} = 200\ A\,\mathrm{GeV}$

Thorsten Renk [a,*], Steffen A. Bass [a,b], Dinesh K. Srivastava [c]

[a] *Department of Physics, Duke University, Durham, NC 27708-0305, USA*
[b] *RIKEN BNL Research Center, Brookhaven National Laboratory, Upton, NY 11973, USA*
[c] *Variable Energy Cyclotron Centre, 1/AF Bidhan Nagar, Kolkata 700 064, India*

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Abstract

We study the role played by the Landau–Pomeranchuk–Migdal (LPM) effect in relativistic collisions of hadrons and heavy nuclei, within a parton cascade model. We find that the LPM effect strongly affects the gluon multiplication due to radiation and considerably alters the spacetime evolution of the dynamics of the collision. It ensures a multiplicity distribution of hadrons in agreement with the experimental proton–proton data. We study the production of single photons in relativistic heavy ion collisions and find that the inclusion of LPM suppression leads to a reduction in the single photon yield at small and intermediate transverse momenta. The parton cascade calculation of the single photon yield including the LPM effect is shown to be in good agreement with the recent PHENIX data taken at the Relativistic Heavy-Ion Collider.

1. Introduction

Collisions of heavy nuclei at relativistic energies are expected to lead to formation of a deconfined phase of strongly interacting nuclear matter, often referred to as a quark–gluon plasma (QGP). Recent data from the Relativistic Heavy-Ion Collider (RHIC) at Brookhaven Lab have provided strong indications for the existence of a transient QGP—among the most exciting findings are strong (hydrodynamic) collective flow [1–6], the suppression of high-p_T particles [7–10] and evidence for parton recombination as hadronization mechanism at intermediate transverse momenta [11–15].

A variety of theoretical models has been formulated to describe the observed phenomena, e.g., fluid dynamical models, perturbative QCD scattering models, as well as models based on parton saturation and statistical approaches. Although these models, which all contain adjustable parameters, have been fairly successful within their respective domains of anticipated applicability, they all have certain limitations. For example, fluid dynamics cannot address the transport phenomena occurring prior to local equilibration of the produced matter, and it must fail above a certain, though unknown, value of p_T. Perturbative parton scattering models fail to describe the physics of equilibration and the evolution of collective flow.

On quite general grounds, one expects that there is an intermediate regime between ultra-hard processes for which the dynamics resembles a superposition of essentially independent nucleon collisions and soft processes which can be described by fluid dynamics. In this regime, semi-hard rescattering of partons produced in primary hard collisions is important, but still perturbatively calculable.

This intermediate regime is addressed by the parton cascade model (PCM), albeit in a semi-classical manner. The parton cascade model [16] was proposed to provide a detailed description of the temporal evolution of nuclear collisions at high energy, from the onset of hard interactions among the partons of the colliding nuclei up to the moment of hadronization. The PCM is based on a relativistic Boltzmann equation for the time evolution of the parton density in phase space due to perturbative QCD interactions including scattering and radiation in the leading logarithmic approximation.

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[*] Corresponding author.
E-mail address: trenk@phys.jyu.fi (T. Renk).

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In the present Letter we report on consequences of Landau–Pomeranchuk–Migdal (LPM) effect on the evolution of the collision within the framework of the PCM approach.

The LPM effect was first derived for QED [17] and describes the destructive interference between bremsstrahlung amplitudes in the case of multiple scattering of the radiating particle. It in effect interpolates between the Bethe–Heitler and factorization limits for the radiation spectrum of a charged particle undergoing multiple scatterings in a medium. The Bethe–Heitler limit is obtained when the separation between individual scattering centers becomes sufficiently large so that the radiation off these centers can be calculated as an incoherent sum of radiation spectra resulting from the individual small-angle scatterings. In the factorization limit the individual scattering centers sit too close together to be resolved by the emitted photon. The observed radiation then factorizes into a product of a single scattering radiation spectrum for the sum of the momentum transfers obtained in all successive small angle scatterings of the emitter and its elastic cross section for this momentum transfer accumulated over these small-angle scatterings. In the regime between those two limiting cases, the LPM effect describes the suppression of radiation relative to the Bethe–Heitler limit in regimes where the radiation formation time is long compared to the mean free path of the emitter and thus destructive interference between the bremsstrahlung amplitudes becomes important. The QED LPM effect has recently been verified by experiments at SLAC [18].

Calculations have shown that the QCD analogue of the LPM effect (it differs from QED due to the non-Abelian nature of QCD) also plays an important role, in particular for estimating the energy loss dE/dx of an energetic parton traversing a dense QCD medium [19–22]. However, all these calculations have focused on limiting cases of infinitely many or very few ($N < 3$) rescatterings of the parton and have either been performed for a static medium or have utilized only very schematic scenarios (e.g., boost-invariant longitudinal expansion) for the evolution of the partonic medium created in a heavy-ion collision.

The LPM effect is expected to limit the growth of parton multiplicity in the dense spacetime-regions of the scattering system. For a perturbative framework like the PCM, this implies that the sensitivity to a soft cut-off scale parameter for particle production is greatly reduced.

We note that other hard scattering models like PYTHIA [23], the DPM [24] and HIJING [25] contain only momentum space physics, they do not include any attempt to address spacetime dynamics. It is therefore difficult to investigate questions dealing with spacetime dynamics (like the LPM effect) within this class of models. All hard scattering models contain parameters associated with, e.g., the separation of hard and soft scales—these parameters are adjusted such that the model describes the p–p data well (in the above cases without rescattering)—but the fact that this is possible for p–p collisions should not be taken as an indication that there is no rescattering or that it would not be important. In [26,27] is has been found that high multiplicity events correspond to underlying events with a high number of hard collisions—it is not unreasonable to suspect that such a high number in the small volume given by the p–p overlap will lead to rescattering (and LPM suppression). In fact, it is the simultaneous description of hard p–p (where rescattering may be observed but is not dominant) and hard A–A physics (where rescattering is expected to be an integral part of the dynamics) with the same underlying set of parameters which is the decisive test for the importance of rescattering and LPM suppression, and we aim to investigate this question in the following.

We describe a schematic implementation of the LPM suppression of gluon and photon radiation into a microscopic transport model, allowing us to study its effect on the non-equilibrium spacetime evolution and reaction dynamics of a heavy-ion collision at RHIC. We show that the same implementation of the LPM suppression is able to describe two sets of data which represent very different conditions in the collision system, i.e., the multiplicity distribution of produced secondary particles in p–p collisions for different $\sqrt{s}$ and photon production in 200 A GeV Au–Au collisions. We select this particular choice of observables because there is no multiplicity distribution data available for heavy-ion collisions since the measured multiplicity is used to determine collision centrality. Direct photon emission is then the cleanest observable measuring the amount of hard collisions taking place in the pQCD rescattering phase.

We argue that this simultaneous agreement demonstrates that we have indeed been able to introduce the physics of LPM suppression into the PCM correctly and that we have a reliable description of the physics in the regime between soft fluid dynamics and the hard pQCD regime where processes scale with the number of binary collisions.

2. The parton cascade model

The fundamental assumption underlying the PCM is that the state of the dense partonic system can be characterized by a set of one-body distribution functions $F_i(x^\mu, p^\alpha)$, where i denotes the flavor index ($i = g, u, \bar{u}, d, \bar{d}, \ldots$) and x^μ, p^α are coordinates in the eight-dimensional phase space. The partons are assumed to be on their mass shell, except before the first scattering. In our numerical implementation, the GRV-HO parametrization [28] is used, and the parton distribution functions are sampled at an initialization scale Q_0^2 ($\approx (p_T^{\min})^2$; see later) to create a discrete set of particles. Partons generally propagate on-shell and along straight-line trajectories between interactions. Before their first collision, partons may have a space-like four-momentum, especially if they are assigned an "intrinsic" transverse momentum.

The time-evolution of the parton distribution is governed by a relativistic Boltzmann equation

$$p^\mu \frac{\partial}{\partial x^\mu} F_i(x, \vec{p}) = C_i[F], \tag{1}$$

where the collision term C_i is a nonlinear functional of the phase-space distribution function. The calculations discussed below include all lowest-order QCD scattering processes between massless quarks and gluons [29], as well as all ($2 \to 2$) processes involving the emission of photons ($qg \to q\gamma$, $\bar{q}g \to$

$\bar{q}\gamma$, $q\bar{q} \to g\gamma$, $q\bar{q} \to \gamma\gamma$). A low momentum-transfer cut-off $p_T^{\min}$ is needed to regularize the infrared divergence of the perturbative parton–parton cross sections. A more detailed description of our implementation is in preparation [30].

We account for the final state radiation [31] following a hard interaction in the parton shower approach. In the leading logarithmic approximation, this picture corresponds to a sequence of nearly collinear $1 \to 2$ branchings: $a \to bc$. Here a is called the mother parton, and b and c are called daughters. Each daughter can branch again, generating a structure with a tree-like topology. We include all such branchings allowed by the strong and electromagnetic interactions. The probability for a parton to branch is given in terms of the variables Q^2 and z. Q^2 is the momentum scale of the branching and z describes the distribution of the energy of the mother parton a among the daughters b and c, such that b takes the fraction z and c the remaining fraction $1 - z$. The differential probability to branch is

$$dP_a = \sum_{b,c} \frac{\alpha_{abc}}{2\pi} P_{a\to bc} \frac{dQ^2}{Q^2}\, dz, \tag{2}$$

where the sum runs over all allowed branchings. The α_{abc} is α_{em} for branchings involving emission of a photon and α_s for the QCD branchings. The splitting kernels $P_{a\to bc}$ are given in [32]. The collinear singularities in the showers are regulated by terminating the branchings when the virtuality of the (time-like) partons drops to μ_0, which we take as 1 GeV. We note that there is no great sensitivity to the detailed choice of μ_0 as soon as the LPM suppression is included (a reduction to $\mu_0 = 0.5$ GeV leads to less than 30% change in parton production) since the LPM effect limits the density of produced partons. In principle, one could take a smaller value for the cut-off μ_0 for a quark fragmenting into a photon [33], but we have not done so as we are only interested in high energy photons here. The soft-gluon interference is included as in [31], namely by selecting the angular ordering of the emitted gluons. An essential difference between emission of a photon and a parton in these processes is that the parton encounters further interactions and contributes to the build-up of the cascade, while the photon almost always (in our approximation, always) leaves the system along with the information about the interaction.

Since the microscopic degrees of freedom of the PCM, quarks and gluons, are treated as quasi-particles, a full quantum implementation of the LPM is beyond the scope of the model. In order to take the main characteristics of the LPM effect into account, we introduce a formation time for the radiated particle,

$$\tau = \frac{\omega}{k_\perp^2}, \tag{3}$$

with ω the energy of the radiated particle and $k_\perp$ its transverse momentum with respect to the emitter. During the formation time, the emitted particles (which we refer to as *shower*) do not interact (and are thus assigned zero cross section). The shower emitter, however, may rescatter and if this occurs within duration of the formation time of the emitted particles, the shower is considered to by suppressed by the LPM effect and is removed from the further evolution of the collision system. Preliminary results of the effects of the LPM suppression on single photon production in the PCM have earlier been reported in [34].

It should be noted that a recent novel implementation of the PCM model [35], based on a stochastic implementation of the collision term (and thus allowing for detailed-balance conserving three-particle collisions), requires a very different modeling of the LPM effect. In that approach the LPM effect is introduced via a lower momentum cut-off for the gluon emission rate, leading to a nearly isotropic angular distribution for inelastic scatterings and thus shorter thermalization times. The effects of this particular implementation on photon radiation and multiplicity scaling at RHIC remain to be investigated.

3. Scaling of multiplicity distributions scaling in *p–p* collisions

In [36] it was suggested that asymptotically the distribution of the multiplicity of produced particles in p–p collisions $\langle n\rangle\sigma_n(s)$ is only a function of $n/\langle n\rangle$, where $\sigma_n(s)$ is the multiplicity distribution for given center of mass energy $\sqrt{s}$ and $\langle n\rangle$ is the mean multiplicity for this s. Thus, the probability distribution $P(n/\langle n\rangle)$ for producing a given fraction of the mean multiplicity would asymptotically be a universal function $\Psi(n/\langle n\rangle)$ independent of $\sqrt{s}$.

In view of this expectation, a large body of data has been accumulated on the multiplicity distribution of hadrons in p–p collisions at several energies, and a description of these have remained an important check on the models of hadronic interactions. Deviations from this (KNO) scaling have also been studied extensively and are most clearly seen in events having high multiplicity at higher center of mass energies (see, e.g., [38]). The high multiplicity events in p–p collisions necessarily involve increased multiple scatterings and gluon multiplications in a small spacetime volume, when studied within models involving scattering of partons (see, e.g., Refs. [26,27]). Thus they provide the most easily tractable arena to study the consequences of the LPM effect.

The PCM should reproduce these multiplicity distributions in order to be reliable. However, there is one important caveat when comparing with data: the PCM does not include hadronization, thus numbers of produced partons in the PCM have to be compared with measured hadron numbers. In the following, we make the assumption that the number of partons produced in a collision scales with the number of measured hadrons, i.e., $N_{\text{part}} \propto N_{\text{had}}$. This assumption has often been made in PCM studies (see Ref. [39]).

Under this assumption, we observe that the PCM without inclusion of the LPM suppression leads to a scaling of the multiplicity distribution in $n/\langle n\rangle$ (Fig. 1, upper panel). However the 'universal function' Ψ^{PCM} is quite different from the measured Ψ^{data}. In particular, there is a large probability to produce a high multiplicity. We note that in the region under investigation, i.e., $\sqrt{s} < 200$ GeV, scaling violations are small and KNO scaling is fulfilled within experimental errors, i.e., the data show a 'universal function'.

Adjusting for the fact that we compare with non-single-diffractive events, we remove the events with zero particle

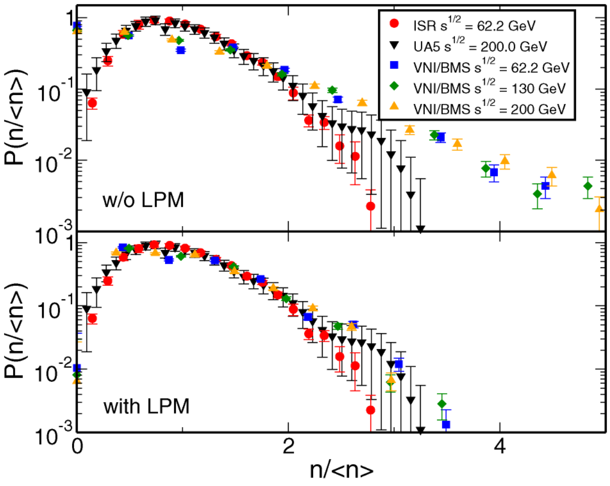

Fig. 1. Upper panel: experimental data from [37,38] and PCM results without LPM reduction for different $\sqrt{s}$ normalized as probability P to find multiplicity n, plotted as a function of the KNO scaling variable $n/\langle n\rangle$. Lower panel: as above, but PCM results include the LPM suppression and non-single diffractive events only.

production from the PCM event sample. Taking the effects of LPM suppression into account, the resulting scaling function Ψ^{PCM} is much closer to Ψ^{data} and it is well conceivable that hadronization can account for the remaining differences. This result gives some confidence that the implementation of the LPM effect is done in a reasonable fashion. We have also confirmed that the PCM (including the LPM effect) is in agreement with the experimental multiplicity distribution for $\sqrt{s} = 900$ GeV [38] with the same level of accuracy as in the comparison shown above.

4. LPM dynamics in Au–Au collisions

We now apply the same implementation of the LPM suppression to Au–Au collisions where a correct treatment of the suppression is even more important due to the higher parton density of the system. As a reference, we investigate Au–Au collisions at 200 GeV/nucleon as realized at RHIC.

In Fig. 2 we compare the collision rate as a function of time for the standard scenario (without LPM suppression) and the one including the LPM effect. In both cases, the collision rate peaks strongly around maximum overlap of the two nuclei at $t = 0$ and then decreases rapidly as subsequent expansion dilutes the system. The LPM effect strongly limits particle production in this high density peak, leading to a collision rate which is almost an order of magnitude lower. Since particle production in the PCM proceeds by branching processes which create soft partons, the result implies that the parton spectra remain harder if the LPM suppression is taken into account.

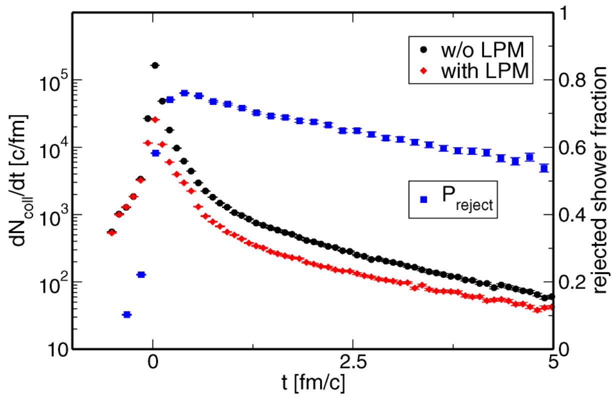

Fig. 2. Left axis: the collision rate in the PCM as a function of time for a scenario with (red) and without (black) LPM suppression. Right axis: the fraction of rejected showers as a function of time (blue). (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

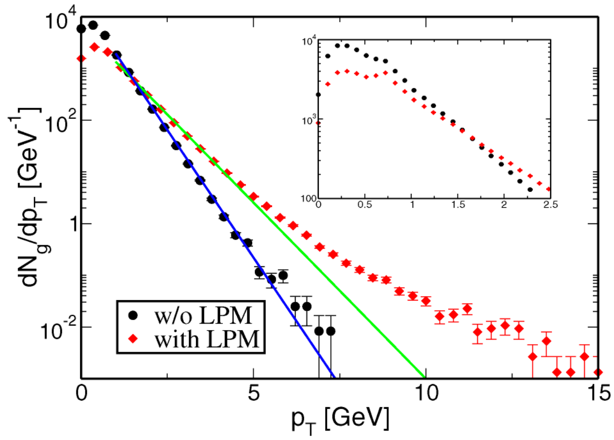

Fig. 3. Transverse mass p_T spectra of gluons for a scenario with (red) and without (black) LPM suppression as obtained from the PCM. The insert shows the low p_T region in greater detail. Shown are also lines indicating exponential fits to the region $1 < p_T < 4$ GeV of the spectra, corresponding to apparent temperatures T^* of 0.64 (0.44) GeV respectively. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

In addition we show the fraction P_{reject} of rejected showers as a function of time. Since in the PCM implementation of the LPM effect the decision about shower rejection is made after a formation time $\tau = \omega/k_\perp^2$, the maximum of the shower rejection does not coincide with the peak in the collision rate but is delayed. The result indicates that the fraction of rejected showers is large throughout the whole evolution phase in which a perturbative interaction picture is expected to be relevant.

This scenario is essentially confirmed by a direct comparison of gluon transverse momentum spectra in Fig. 3. The much higher collision rates in the standard scenario lead to an exponential spectrum with an apparent temperature T^* (determined by a fit $\sim \exp[-p_T/T^*]$ in the range $1 < p_T < 4$ GeV) of 0.44 GeV. In contrast, including the LPM suppression leads to less spectral cooling (the apparent temperature is 0.64 GeV in the region where the spectrum is exponential) and the power-law tail remains clearly visible for $p_T > 5$ GeV (note that T^* is not a real temperature as for example the longitudinal momentum spectrum looks very different from the transverse one and dN/dE does not follow the thermal distribution). On the other hand, focusing on the low p_T region reveals that including the LPM suppression leads to a factor ~ 2.5 reduction in the yield below 0.7 GeV (the $p_{T,\mathrm{min}}$ cut-off).

5. Photon production in Au–Au collisions

Ideally we would like to study a multiplicity scaling plot of Au–Au collisions to directly compare to our p–p collision results. However, experimentally produced multiplicity is used to determine the centrality class of the collision, hence a multiplicity distribution for central Au–Au collisions is not available and we have to rely on a different probe sensitive to the number of hard collisions. Due to the smallness of the electromagnetic coupling, photons produced in a heavy-ion collision can escape the interaction region essentially unaltered [40] and are therefore an excellent and reliable probe of the evolution of the partonic cascade in nuclear collisions. In fact, one can argue that the photons confirm the presence of parton cascading processes after the initial primary parton–parton collisions. If we would attempt to describe p–p collisions without rescattering processes and to carry the same description over to heavy-ion collisions the resulting photon yield is reduced by about a factor of ~ 3–4 (cf. also the discussion in [34]).

Photons are produced in the PCM from Compton ($qg \to q\gamma$), annihilation ($q\bar{q} \to g\gamma$), and bremsstrahlung ($q^\star \to q\gamma$) processes. These are analogous to the processes governing the energy loss of energetic partons, where gluons are emitted instead of photons. As in [34] we investigate the photon production during the hard initial stage of Au–Au collisions, focusing here on the effect of the LPM suppression on single photon emission. We find a sizable reduction of photon production in particular at midrapidity (where experimental measurements have been made) as compared to [34] due to LPM suppression.

Fig. 4 shows the p_T spectrum of single photons calculated with and without LPM suppression as compared to the data obtained by the PHENIX Collaboration [41] for the 0–10% centrality class.

We clearly observe that without the inclusion of the LPM suppression the PCM overpredicts the data in the region $2 < p_T < 3.5$ GeV whereas taking the LPM effect into account leads to a fair description of the data. A thermal contribution to the photon yield, calculated using a fireball model [42,43] which reproduces HBT correlations for pions, or using a boost-invariant hydrodynamics with transverse expansion [44], with state of the art rates for thermal emission of photons from quark and hadronic matter, seems to improve the agreement with data in the low p_T region, although the errors are large here. Note that without the presence of the LPM effect, frequent soft scattering and branching processes generate a lot of photons below 0.7 GeV ($= p_T^{\min}$). These processes are strongly suppressed by the LPM effect, explaining the difference between a rise of the low p_T spectrum (without LPM) and a drop (including LPM) in the PCM results.

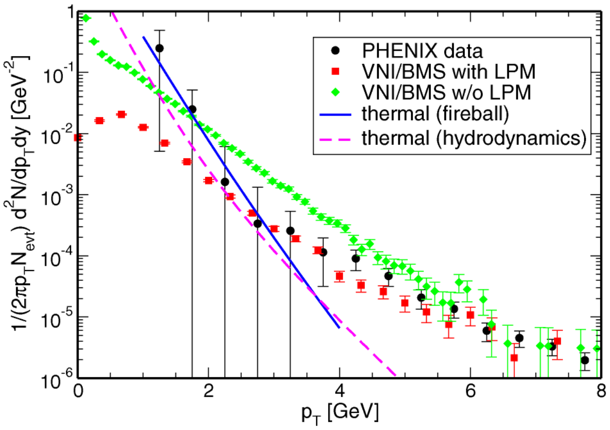

Fig. 4. Photon transverse momentum spectra at midrapidity, as measured by the PHENIX Collaboration [41] and as calculated in the PCM with (red) and without (green) inclusion of the LPM effect. Shown is also a calculation of the thermal contribution to the spectrum using a fireball model of expansion based on a fit to hadronic data (solid blue line), [42,43], and a hydrodynamic calculation (dashed magenta) [44]. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

6. Conclusions

We have studied the role played by the Landau–Pomeranchuk–Migdal (LPM) effect in relativistic collisions of hadrons and heavy nuclei, within the framework of the parton cascade model. We find that the LPM effect strongly affects the gluon multiplication due to radiation and considerably alters the spacetime evolution of the dynamics of the collision. In particular, it restricts the growth of multiplicity due to soft parton production considerably and strongly reduces the sensitivity of parton production to the detailed choice of a soft cut-off parameter μ_0. It ensures a multiplicity distribution of hadrons in agreement with the experimental data in proton–proton reactions. Furthermore, we find that the production of single photons in relativistic heavy ion collisions is strongly affected–the inclusion of the LPM suppression leads to a depletion of single photons at low and intermediate transverse momenta up to 4 GeV/c and brings the PCM calculation into good agreement with the recent PHENIX data taken at the Relativistic Heavy-Ion Collider. The success in the reproduction of these two different sets of data is far from trivial as they represent not only very different observables (bulk production vs. rare process) but also vastly different $\sqrt{s}$ and parton densities. This success gives some confidence that we have a useful description of the regime in which collisions are still perturbatively calculable but multiple rescattering is important.

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