

Energy–Energy Correlator for jet production in pp and pA collisions

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In this Letter, we study the collinear limit of the Energy–Energy Correlator (EEC) in single-inclusive jet production in proton–proton (pp) and proton–nucleus (pA) collisions. We introduce a non-perturbative model that allows us to describe the EEC in the entire angular region of the current experiments. Our results for pp collisions show excellent agreement with CMS and ALICE data over a wide range of jet transverse momenta. For pA collisions, we include modifications from the nuclear medium, and our predictions align well with the trends observed in recent ALICE measurements.

Introduction. Hadronic jets—energetic, collimated cascades of QCD particles produced in high-energy scattering experiments—are among the most powerful tools for probing the nature and properties of QCD. In events with low final-state particle multiplicity, such as those typically observed in ee and pp collisions, jets enable precise tests of the soft and collinear dynamics of QCD, see e.g. [1–3]. In contrast, in collisional systems such as eA , pA , and AA , the jet radiative cascade gets modified by the underlying QCD medium, whether it be the cold nuclear matter characterized by a highly occupied gluonic state [4, 5], or the hot quark–gluon plasma (QGP) [6]. By comparing these modified jets to those produced in ee or pp collisions, one can investigate the medium-induced imprints on the jet cascade, providing a unique window into the many-body properties of QCD.

One of the key challenges in describing the vacuum fragmentation pattern of jets and extracting medium-induced modifications to jets lies in identifying observables that allow for both a theoretically consistent description and experimental accessibility across a broad kinematic region. Recently, it has been argued that correlations between asymptotic fluxes of certain currents can satisfy these requirements, see e.g. [7–17]. Focusing on the widely studied case of Energy–Energy Correlators (EECs) [18–21], in vacuum they allow, for example, to directly access the anomalous scaling of leading-twist QCD operators at high precision [8, 9, 20, 22] and cleanly separate the relevant scales inside jets, sharply splitting perturbatively and non-perturbatively dominated sectors [8, 17, 22]. Extending this program to eA , pA , and AA collisions, EECs have the potential to directly probe the scales of the underlying matter and resolve its spatial structure [12, 23–35].

Despite this promise, the theoretical control of EECs in complex collisional systems, particularly AA collisions, remains limited. Current studies are mostly constrained to (partial) leading-order perturbative calculations [24, 25, 36] and numerical model simulations [25, 34]. Even

in cleaner environments such as ep or pp collisions, the description of EECs in the non-perturbative sector is still poorly understood.

This non-perturbative region of EECs has recently garnered a lot of interest due to the emergent scaling behavior, indicative of the hadronization transition’s universality [15–17, 22, 37]. In the back-to-back limit of EECs where the transverse-momentum-dependent (TMD) factorization is valid, the non-perturbative region can be described using non-perturbative Sudakov factors, related to the transverse momentum dependence of fragmentation, and the non-perturbative Collins–Soper (CS) kernel required for a resummation of large logarithms in rapidity [11, 38–41]. In the collinear limit, the CS resummation is not needed, but the physics related to the non-perturbative fragmentation is expected to remain the same, see e.g. [42]. Exploring the universality of these non-perturbative effects is then of high interest due to its potential in describing a vast range of observables across different experimental environments.

To develop a comprehensive understanding of jet substructure and medium-induced modifications from EECs, it is essential to have theoretical predictions that encompass both the perturbative and the non-perturbative sectors of pp collisions, while also incorporating medium-induced modifications for systems interpolating between pp and AA events. A prime candidate for this is pA collisions where high-energy jets can still be produced and fragment in the presence of cold nuclear matter [43–51]. This constitutes a *cleaner* environment when compared to the QGP, allowing us to directly probe the medium-induced modifications to EECs inside jets.

In this Letter, we present a formalism to describe the collinear limit of the EEC distribution in pp and pA collisions. Our framework accounts, at the same time, for both the non-perturbative regime, where a parametrization inspired by TMD factorization is implemented [11, 52], and the perturbative sector, where we employ the leading-logarithmic (LL) evolution of the

vacuum cascade as well as the medium-induced modifications to the partonic branchings in pA events. This holistic approach allows for an efficient description of the entire EEC region and can be extended to other collisional systems. As we shall show, this permits a first direct observation of purely medium-induced modifications to the jet cascade in pA collisions using EECs.

Proton–proton collisions. The collinear limit of the EEC distribution satisfies the factorization formula [8, 9]:

$$\frac{d\Sigma^{pp}}{dp_T dy dR_L} = \int_0^1 dx x^2 \frac{dJ(x, p_T, R_L)}{dR_L} \cdot H(x, y, p_T), \quad (1)$$

where we made implicit the dependence on the factorization scale μ , y denotes the rapidity, and p_T is the transverse momentum of the jet. This factorization formula separates the production of the outgoing states into a hard function H , which does not depend on the angular separation between the outgoing states R_L , and an EEC jet function J , which incorporates the collinear fragmentation inside the jet. Restraining our discussion to the LL accuracy, the convolution in Eq. (1) trivializes to setting $x = 1$, and the dot product is taken in flavor space, i.e. $J = \{J_q, J_g\}$, and respectively for H .

The hard function can be obtained from collinear factorization [53–56], where the cross section to extract a parton c from the initial hard scattering reads

$$\frac{d\sigma_c}{dp_T dy} = \sum_{a,b} f_{a/p}(x_a, \mu) \otimes f_{b/p}(x_b, \mu) \otimes \hat{\sigma}_{a+b \rightarrow c}. \quad (2)$$

Here \otimes denotes a convolution over the parton momentum fraction, $\hat{\sigma}_{ab \rightarrow c}$ is the hard partonic cross section to produce a parton c with a transverse momentum p_T and rapidity y , and $f_{a/p}$ is the parton distribution function (PDF) associated to the parton a coming from the proton. The hard function can then be defined from Eq. (2) as $H(y, p_T, \mu) = \{d\sigma_q/(\sigma dp_T dy), d\sigma_g/(\sigma dp_T dy)\}$ [57–60], which is normalized by the jet cross section σ as in the definition of the EEC. We choose the factorization scale $\mu = p_T$.

Moving to the jet function, we first note that renormalization group (RG) consistency implies that, at the LL accuracy, J satisfies a multiplicative renormalization formula, which results in the differential jet function [9]

$$\frac{dJ}{dR_L} = \frac{\alpha_s(R_L p_T)}{\pi R_L} (1, 1) \cdot \left(\frac{\alpha_s(\mu)}{\alpha_s(R_L p_T)} \right)^{\frac{\gamma(3)}{\beta_0}} \cdot \gamma(3), \quad (3)$$

where $\beta_0 = (11C_A - 2n_f)/3$ and $\gamma(j) \equiv -\int_0^1 dz z^{j-1} \hat{P}(z)$ is the anomalous dimension matrix, with \hat{P} being the regularized splitting function, such that only the spin $j = 3$ contributes to the EEC at the leading power. As a result, Eq. (3) defines the perturbative vacuum component of the EEC jet function.

So far, we have discussed the perturbative sector where $R_L p_T \gg \Lambda_{\text{QCD}}$. At lower scales, the EEC becomes highly sensitive to non-perturbative dynamics in the fragmentation process. The underlying physics of fragmentation leading to the modification of the momentum distribution should be the same as in the back-to-back limit, and thus we model it using a form inspired by the non-perturbative Sudakov factor in TMD factorization [11, 42]:

$$j_{\text{np}}(b) = \exp(-a_0 b), \quad (4)$$

where a_0 is a *free* parameter accounting for non-perturbative effects. We note that the parameter a_0 should differ for quarks and gluons due to their different color charge, but in this work we shall not consider the separation in flavor. In the momentum space, we can interpret j_{np} as a probability distribution for the produced hadrons to get additional transverse momentum in the fragmentation, and the condition $j_{\text{np}}(b = 0) = 1$ guarantees that the distribution is correctly normalized. Furthermore, we note that this non-perturbative element inherits the universal behavior of the hadronization transition and can thus be applied to extractions of the (vacuum) jet EEC in other processes; we discuss the medium-induced modifications below. To incorporate this non-perturbative element, it is convenient to work in the position space where the convolution is multiplicative:

$$\frac{d\Sigma^{pp}}{dp_T dy dR_L} = R_L p_T^2 \int_0^\infty db b J_0(R_L p_T b) j_{\text{np}}(b) \tilde{\Sigma}(b), \quad (5)$$

$$\tilde{\Sigma}(b) = (1, 1) \cdot \left(\frac{\alpha_s(R p_T)}{\alpha_s(\mu b_*)} \right)^{\frac{\gamma(3)}{\beta_0}} \cdot H. \quad (6)$$

This relation follows from evolving the jet function in the position space where the initial scale $R_L p_T$ is replaced by $\mu b_* = 2e^{-\gamma_E}/b_*$, see supplementary material. This choice minimizes large logarithms that appear at higher orders, and we use the standard b_* -prescription $b_* = b/\sqrt{1 + b^2/b_{\text{max}}^2}$ with $b_{\text{max}} = 2e^{-\gamma_E} \text{GeV}^{-1}$ [61, 62]. Eq. (5) resums the same logarithms as the momentum-space expression in Eq. (3). We take into account the finite jet radius $R = 0.4$, consistent with the experimental data, by evolving the jet function up to $R p_T$ [63].

The numerical results for the EEC in pp collisions are shown in Fig. 1, with the parameter a_0 characterizing non-perturbative physics chosen as $a_0 = 3.8 \text{ GeV}$ for CMS and $a_0 = 2.5 \text{ GeV}$ for ALICE, fitted to the experimental data presented in [16] and [17], respectively. We use the CT18ANLO fit [64] for the proton PDF, and we also incorporate a K -factor of $K \approx 0.90$ for CMS and $K \approx 0.75$ for ALICE to account for higher-order corrections beyond the LL approximation in our formalism. While the non-perturbative part is supposed to be universal, we use different parameter values to describe the CMS and ALICE data. The reason for this is that

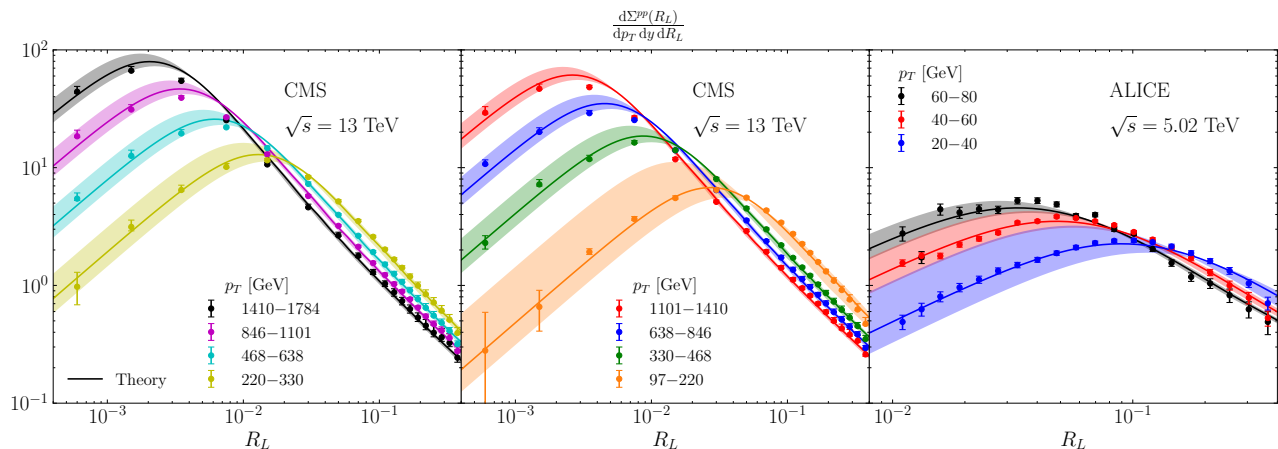


FIG. 1. Differential EEC distribution in pp collisions as a function of the angle R_L compared to the experimental data from CMS [16] and ALICE [17]. The theory prediction at the average p_T of the bin is shown with solid lines, while the shaded band corresponds to the prediction with the minimum and maximum p_T in each bin.

the ALICE data is only for charged jets while the CMS data is for inclusive jets, and thus the non-perturbative parameters should differ.

A direct analysis of Fig. 1 allows us to conclude that our description, accommodating the LL perturbative prediction at large angles and a non-perturbative jet function at small angles, can describe the data quite precisely in the entire angular domain and for a wide range of different p_T values. We note that while the large-angle region is fully determined from perturbative considerations, the non-perturbative sector for all of the curves in each experiment is described by the same a_0 parameter, in agreement with the universal character of j_{np} . Most crucially, we capture the location of the distribution's peak for all values of p_T , indicating a good theoretical description of the transition between the two sectors.

Proton-nucleus collisions. Having introduced a description of the EEC in pp collisions in the full angular domain, we extend our construction to collisional systems where a QCD background medium is present. The presence of the medium can, in general, induce modifications to both the hard and the EEC jet functions. In pA collisions, the hard function should be updated to account for the target nucleus. To that end we replace one of the PDFs in Eq. (2) by a nuclear PDF (nPDF), i.e. $f_{b/p}(x_b, \mu) \rightarrow f_{b/A}(x_b, \mu)$.

The jet function receives medium-induced modifications to the perturbative and non-perturbative sectors. For the former, the medium correction enters as a higher-twist/power correction to the leading perturbative QCD result. Thus, it cannot change the RG evolution at the current accuracy, and the corrections can be entirely incorporated into the fixed-order jet function which now reads $J^{pA} = J^{pp} + J^{\text{med}}$, with

$$\frac{dJ^{\text{med}}}{dR_L} = \frac{\alpha_s(R_L p_T)}{\pi R_L}$$

$$\times \int_0^1 dx x(1-x) P^{\text{vac}}(x) F_{\text{med}}(p_T, R_L, x), \quad (7)$$

where the medium modification (matrix) factor F_{med} depends on the jet quenching parameter \hat{q} and the effective path length in the medium L . Here P^{vac} denotes the appropriate vacuum splitting kernel. In the large- N_c limit, F_{med} can be explicitly computed in a semi-classical expansion [65, 66]; we provide the explicit formulas in the supplementary material.

The non-perturbative sector of J^{pA} is also modified compared to pp collisions. Physically, these corrections are driven by the multiple scattering of the jet partons on the nucleus, which results in the broadening of the underlying momentum distribution. Thus, at small angles we update Eq. (4) to the form

$$j_{\text{np}}(b) = \exp(-a_0 b - a_1 b^2), \quad (8)$$

where a_0 is the same as in the vacuum case and the parameter a_1 accounts for the medium-induced diffusive contribution to the hadronization of the jet partons. Note that a_1 captures the relative transverse momentum broadening of a particle pair inside the jet, differing, in general, from the standard $\hat{q}L$ scaling for a single parton. Finally, we also note that, as in AA collisions, one should account for energy-loss effects which can affect the non-perturbative sector. However, since in pA collisions these are expected to be negligible due to the smallness of \hat{q} and L , we shall neglect them in what follows.

We consider a typical value for the jet quenching parameter in cold nuclear matter, $\hat{q} = 0.02 \text{ GeV}^2/\text{fm}$ [67, 68], and take the path length to be $L = 3 \text{ fm}$, while the non-perturbative medium-dependent parameter is set at $a_1 = 0.25 \text{ GeV}^2$. The K -factors, the proton PDF and the a_0 parameter are the same as in our pp analysis, while we use the EPPS21 fit [69] for the nPDF.

The numerical results for the ratio between the differential EEC in $p\text{Pb}$ with respect to pp collisions as a function of the opening angle, defined as $R_{pA} = \frac{d\Sigma^{pA}(R_L)}{dp_T dy dR_L} / \frac{d\Sigma^{pp}(R_L)}{dp_T dy dR_L}$, are shown in Fig. 2, along with the fit to the experimental data in the $p_T \in (20, 40)$ GeV bin provided by the ALICE collaboration [70] (dashed black line). We first note that the effect of the nPDF (solid green line) is almost negligible for the whole R_L range. This is expected since the EEC is normalized by the jet cross section, and the effect of nPDF mostly cancels out in the ratio. We then present two predictions, one including only the medium modification to the non-perturbative sector (solid red line) and one including all the medium effects (solid blue line). As is evident when comparing to the experimental fit, our modification to j_{np} is crucial to capture the data trend for a suppression in R_{pA} . This gives a clear indication of the importance of momentum broadening in the nucleus at small transverse momentum scales. We note that, so far and to our best knowledge, momentum diffusion had only been taken into account in the perturbative region of the EEC.

Comparing the experimental fit to the theory at large R_L , we clearly observe the enhancement effect driven by the medium modification factor F_{med} . Note that at small angles $F_{\text{med}} \approx 0$, and thus it does not affect the non-perturbative region. Moreover, compared to the case of AA collisions where there is a competing effect driven by the soft physics of the bulk [25, 34, 71], in the pA scenario the medium does not significantly contaminate the jet. As a result, our analysis provides a first clear and unambiguous evidence for the emergence of medium-induced jet modifications in the EEC distribution. We further note that the approximations used to compute F_{med} in this work tend to overestimate the enhancement, and thus improved theoretical descriptions would most likely drive the model curve closer to the experimental fit, see discussion in [24]. We also provide a prediction for jets with $p_T \in (60, 80)$ GeV in the lower panel of Fig. 2 to illustrate the p_T -dependence of the nuclear modification.

Conclusions. In this Letter, we have presented a theoretical framework for the differential EEC distribution in the collinear limit, encompassing both pp and pA collisions. Our model incorporates the leading-logarithmic form of the observable in the perturbative sector and employs a TMD-inspired parametrization to describe the non-perturbative regime. By fitting the available CMS and ALICE pp data, we have achieved excellent agreement with experiment, successfully describing the EEC at both small and large angles and accurately predicting the location of the distribution's peak. For pA collisions, our formalism incorporates nuclear medium effects, enabling a unified description across both the non-perturbative and the perturbative regions.

In the small-angle region, we have highlighted the critical role played by medium-induced broadening, essen-

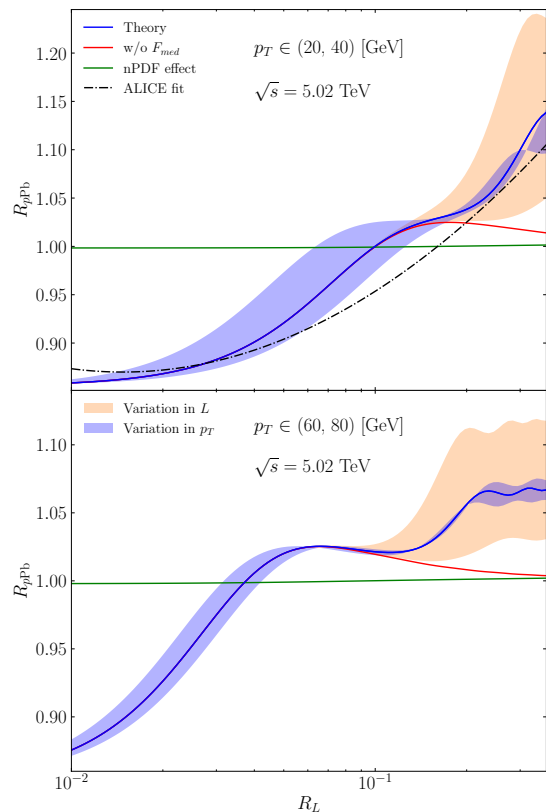


FIG. 2. Ratio of the differential EEC $p\text{Pb}$ distribution with respect to pp collisions, as a function of R_L , compared to the experimental fit provided by ALICE [70] (dashed black line): $R_{p\text{Pb}}^{\text{ALICE}} = 0.87 + 0.121(\log_{10}(R_L) + 1.831)^2$. The theory prediction, at the average p_T of the bin, is shown by the blue solid line. The ratio with $F_{\text{med}} = 0$ (solid red line), and the sole effect of introducing the nPDF in the vacuum distribution (solid green line) are also shown. The blue shaded band corresponds to the prediction using the minimum and maximum p_T of the bin, while the orange band corresponds to modifying the length of the medium $L \in (2, 4)$ fm.

tial to accurately capture the hadronization transition. At large angles, we have included the leading medium-induced power corrections, demonstrating that these lead to a significant enhancement compared to pp data. Contrasting our theoretical predictions with ALICE's fit to experimental data [70], we have shown that our framework provides a precise and comprehensive description of the measurements.

This work establishes a solid foundation for a future detailed and thorough understanding of the jet EEC across a diverse set of experiments, opening exciting opportunities for future research.

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SUPPLEMENTARY MATERIAL

Evolution of the jet function in position space

In this appendix, we describe the evolution of the jet in the position space instead of the momentum space. We define the position-space jet function $j(b)$ as the Fourier transform of the differential jet function $dJ/d^2\vec{k}_T$,

$$j(b) = \int d^2\vec{k}_T \frac{dJ}{d^2\vec{k}_T} e^{i\vec{b}\cdot\vec{k}_T}, \quad (9)$$

where $k_T = p_T R_L$ and

$$\frac{dJ}{d^2\vec{k}_T} = \frac{1}{2\pi R_L p_T^2} \frac{dJ}{dR_L} = \delta^2(\vec{k}_T) \cdot (1, 1) \quad (10)$$

at leading order. This corresponds to $j(b) = (1, 1)$ in the position space.

The evolution equation for the jet function resumming the large logarithms is not modified when going from the momentum to the position space, and it reads:

$$\frac{dj(b, \mu)}{d \ln \mu^2} = -\frac{\alpha_s(\mu)}{4\pi} j(b, \mu) \cdot \gamma(3). \quad (11)$$

Keeping only the LO β -function in the running of the coupling,

$$\frac{d\alpha_s(\mu)}{d \ln \mu^2} = -\frac{1}{4\pi} \alpha_s^2(\mu) \beta_0, \quad (12)$$

the solution to Eq. (11) can be written as

$$j(b, \mu) = j(b, \mu_0) \cdot \left(\frac{\alpha_s(\mu)}{\alpha_s(\mu_0)} \right)^{\frac{\gamma(3)}{\beta_0}}. \quad (13)$$

Evolving to the final scale $\mu = Rp_T$ and taking the initial condition as $j(b, \mu_{b_*}) = (1, 1)$, we then get

$$j(b, Rp_T) = (1, 1) \cdot \left(\frac{\alpha_s(Rp_T)}{\alpha_s(\mu_{b_*})} \right)^{\frac{\gamma(3)}{\beta_0}}. \quad (14)$$

This can be combined with the hard function to define the position-space EEC

$$\tilde{\Sigma}(b) = j(b, Rp_T) \cdot H, \quad (15)$$

which corresponds to Eq. (6) in the main article.

Medium modification factors

In this appendix, we provide the explicit form of both the vacuum splitting functions and the modification factors due to the propagation inside a medium of length L for all the relevant channels: $q \rightarrow qq$, $g \rightarrow gg$ and $g \rightarrow q\bar{q}$. The unregularized vacuum splitting functions read

$$\begin{aligned} P_{qq}^{\text{vac}}(x) &= C_F \frac{1+x^2}{1-x}, & P_{gq}^{\text{vac}}(x) &= P_{qq}^{\text{vac}}(1-x), \\ P_{qg}^{\text{vac}}(x) &= 2n_f T_R (x^2 + (1-x)^2), & P_{gg}^{\text{vac}}(x) &= 2C_A \left(\frac{x}{1-x} + \frac{1-x}{x} + x(1-x) \right). \end{aligned} \quad (16)$$

By following the procedure developed in [65, 66], see also [72, 73], we obtain the medium modifications factors. For any flavor channel, they can be written as

$$F_{\text{med}}(x, R_L) = 2 \int_0^L \frac{dt_1}{t_f} \left[\int_{t_1}^L \frac{dt_2}{t_f} \cos\left(\frac{t_2 - t_1}{t_f}\right) \mathcal{C}^{(4)}(L, t_2) \mathcal{C}^{(3)}(t_2, t_1) - \sin\left(\frac{L - t_1}{t_f}\right) \mathcal{C}^{(3)}(L, t_1) \right], \quad (17)$$

with $\mathcal{C}^{(n)}(t_f, t_i)$ the n -particle correlator of light-like Wilson lines in the light-cone time interval (t_f, t_i) , which can be computed in the large- N_c limit by using the harmonic oscillator approximation. Note that to obtain Eq. (17) we have resorted to the so-called *tilted* Wilson line approximation, see e.g. [74].

1. $q \rightarrow qg$: The quark–gluon splitting modification factor is deduced in the aforementioned works and reads

$$F_{\text{med}}^{qq}(x, R_L) = 2 \int_0^L \frac{dt_1}{t_f} \left[\int_{t_1}^L \frac{dt_2}{t_f} \cos\left(\frac{t_2 - t_1}{t_f}\right) e^{-\frac{\hat{q}}{12} R_L^2 (t_2 - t_1)^3 (1 + (1-x)^2)} e^{-\frac{\hat{q}}{4} R_L^2 (L - t_2) (t_2 - t_1)^2 (1 - 2(1-x) + 3(1-x)^2)} \right. \\ \left. \times \left(1 - \frac{\hat{q}}{2} R_L^2 x (1-x) (t_2 - t_1)^2 \int_{t_2}^L ds e^{-\frac{\hat{q}}{12} R_L^2 ((s - t_2)^2 (2s - 3t_1 + t_2) + 6x(1-x)(s - t_2)(t_2 - t_2)^2)} \right) \right. \\ \left. - \sin\left(\frac{L - t_1}{t_f}\right) e^{-\frac{\hat{q}}{12} R_L^2 (L - t_1)^3 (1 + (1-x)^2)} \right], \quad (18)$$

with R_L being the angle between the two prongs, x their energy fraction, p_T the transverse momentum of the jet, and the formation time defined as $t_f = \frac{2}{x(1-x)R_L^2 p_T}$.

2. $g \rightarrow gg$: For this channel, the three-particle correlator is also directly given in the above works:

$$\mathcal{C}_{g \rightarrow gg}^{(3)}(t_2, t_1) = e^{-\frac{\hat{q}}{6} R_L^2 (t_2, t_1)^3 (1 - x + x^2)}. \quad (19)$$

In order to obtain $\mathcal{C}^{(4)}$ for the gluon splitting, the differential equation Eq. (4.12) in [65] must be solved. Doing so, one gets

$$\mathcal{C}_{g \rightarrow gg}^{(4)}(L, t_2) = e^{-\frac{\hat{q}}{4} R_L^2 (L - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)} e^{-\frac{\hat{q}}{12} R_L^2 ((L - t_1)^3 + (L - t_2)^3 - (t_2 - t_1)^3)} \\ \times \left(1 + \frac{\hat{q}}{2} R_L^2 \int_{t_2}^L ds (s - t_1)(s - t_2) e^{\frac{\hat{q}}{12} R_L^2 ((s - t_1)^2 + (s - t_2)^3 - (t_2 - t_1)^3)} e^{-\frac{\hat{q}}{4} R_L^2 (s - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)} \right). \quad (20)$$

Thus, the medium modification factor for the gluon splitting channel can be written as

$$F_{\text{med}}^{gg}(x, R_L) = 2 \int_0^L \frac{dt_1}{t_f} \left[\int_{t_1}^L \frac{dt_2}{t_f} \cos\left(\frac{t_2 - t_1}{t_f}\right) e^{-\frac{\hat{q}}{4} R_L^2 (L - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)} \right. \\ \left. \times e^{-\frac{\hat{q}}{12} R_L^2 ((L - t_1)^3 + (L - t_2)^3 - (t_2 - t_1)^3)} e^{-\frac{\hat{q}}{6} R_L^2 (t_2 - t_1)^3 (1 - x + x^2)} \right. \\ \left. \times \left(1 + \frac{\hat{q}}{2} R_L^2 \int_{t_2}^L ds (s - t_1)(s - t_2) e^{\frac{\hat{q}}{12} R_L^2 ((s - t_1)^2 + (s - t_2)^3 - (t_2 - t_1)^3)} e^{-\frac{\hat{q}}{4} R_L^2 (s - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)} \right) \right. \\ \left. - \sin\left(\frac{L - t_1}{t_f}\right) e^{-\frac{\hat{q}}{6} R_L^2 (L - t_1)^3 (1 - x + x^2)} \right]. \quad (21)$$

3. $g \rightarrow q\bar{q}$: Finally, the medium correction to the pair production channel can be easily computed in the corresponding limits. The three-particle correlator reads

$$\mathcal{C}_{g \rightarrow q\bar{q}}^{(3)}(t_2, t_1) = e^{-\frac{\hat{q}}{12} R_L^2 (t_2 - t_1)^3 (x^2 + (1-x)^2)}, \quad (22)$$

while the four-particle correlator reads

$$\mathcal{C}_{g \rightarrow q\bar{q}}^{(4)}(L, t_2) = e^{-\frac{\hat{q}}{4} R_L^2 (L - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)}. \quad (23)$$

We have noticed a misprint in the equation providing $\mathcal{C}^{(4)}$ in terms of correlators of Wilson lines in [65] for this channel, where the quadrupole and the double dipole are switched, c.f. [75, 76]. With this, the medium modification factor becomes

$$F_{\text{med}}^{qg}(x, R_L) = 2 \int_0^L \frac{dt_1}{t_f} \left[\int_{t_1}^L \frac{dt_2}{t_f} \cos\left(\frac{t_2 - t_1}{t_f}\right) e^{-\frac{\hat{q}}{4} R_L^2 (L - t_2) (t_2 - t_1)^2 (x^2 + (1-x)^2)} e^{-\frac{\hat{q}}{12} R_L^2 (t_2 - t_1)^3 (x^2 + (1-x)^2)} \right. \\ \left. - \sin\left(\frac{L - t_1}{t_f}\right) e^{-\frac{\hat{q}}{12} R_L^2 (L - t_1)^3 (x^2 + (1-x)^2)} \right]. \quad (24)$$