Antenna Showers with Hadronic Initial States

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Abstract

We present an antenna shower formalism including contributions from initial-state partons and corresponding backwards evolution. We give a set of phase-space maps and antenna functions for massless partons which define a complete shower formalism suitable for computing observables with hadronic initial states. We focus on the initial-state components: initial–initial and initial–final antenna configurations. The formalism includes comprehensive possibilities for uncertainty estimates. We report on some preliminary results obtained with an implementation in the \textit{Vincia} antenna-shower framework.

Keywords: parton showers; quantum chromodynamics; hadronic collisions

1. Introduction

Parton-shower algorithms offer a universal and fully exclusive perturbative resummation framework for high-energy processes. In the context of Monte Carlo event generators \cite{1}, they also provide the perturbative input for hadronization models. As such, they are complementary to more inclusive techniques, such as fixed-order calculations (limited to small numbers of hard and well-separated partons) and more inclusive resummation approaches (limited to a fixed set of observables).

Sjöstrand derived the first consistent parton-shower algorithm \cite{2} for so-called “backwards evolution” of initial-state partons a quarter-century ago. The central point is that an initial-state parton defined at a high factorization scale, $Q_F$, can be evolved “backwards”, towards earlier times, to find the parton from which it originated at some low scale, $Q_0 \sim 1$ GeV. During this evolution, which is governed by the Altarelli-Parisi splitting kernels \cite{3} supplemented by PDF ratios (a point which is crucial to the backwards-evolution formalism), initial-state radiation is emitted, which in turn gives rise to its own final-state radiation, and the character of the evolving parton changes, migrating towards successively higher $x$ values and towards the more valence-dominated flavor content at low $Q$.

As an alternative to Altarelli-Parisi evolution, Gustafson and Pettersson proposed a final-state algorithm based on QCD dipoles \cite{4}, which has been implemented in \textit{Ariadne} \cite{5}. There, however, initial-state radiation does not rely on backwards evolution. Instead, it is treated essentially as final-state radiation off dipoles stretched between the hard process and the beam remnants, and thus depends on the non-perturbative makeup of the remnants. Winter and Krauss took a first step towards combining the dipole formalism with backwards evolution (and thus also eliminating the dependence on the remnants) in ref. \cite{6}. Our construction differs in the antenna functions, evolution variables, and recoil strategy. In particular, it differs in the treatment of collinear singularities in initial–final antennæ. We have checked that our antennæ properly reproduce all QCD singularities.

A complementary approach which merges the Lund dipole language with that of fixed-order antenna factorization \cite{7} \cite{8} \cite{9} \cite{10}, is that of \textit{Vincia}...
the term “antenna” rather than “dipole” to avoid ambiguities of historical origins, see e.g., ref. [14]. So far, however, the VINCIA formalism has been applied only to final-state showers. In this paper, we present all the ingredients necessary to construct a consistent initial-state shower based on QCD antennæ. A further important ingredient is comprehensive possibilities for uncertainty estimates, in line with the framework for automated theory uncertainties proposed in ref. [15].

2. Antennæ and Antenna Showers

Throughout this Letter, we use the following notation convention: capital letters for pre-branching (parent) partons and lower-case letters for post-branching (daughter) ones. Also, we use the first letters of the alphabet, a, b, c, ..., for incoming partons and letters starting from i, j, k, ... for outgoing ones. Fig. [1] illustrates these choices for the two basic types of configurations we consider. We will also mark incoming particles with a minus sign in front in antenna functions. We adopt the convention that particle energies are always positive, whether the particle is in the initial or the final state. As a result, \( s_{ij} = (k_i + k_j)^2 \) is always positive.

The key building block for parton showers is the Sudakov factor, which represents the non-emission probability between two values of the evolution scale, see [1] for reviews. In the context of an antenna shower, the Sudakov factor for the branching of one antenna is

\[
\Delta \left( Q^2_{\text{start}}, Q^2_{\text{emit}} \right) = \exp \left[ -\mathcal{A}(Q^2_{\text{start}}, Q^2_{\text{emit}}) \right],
\]

with

\[
\mathcal{A}(Q^2_{\text{start}}, Q^2_{\text{emit}}) = \int_{Q^2_{\text{start}}}^{Q^2_{\text{emit}}} a_c(\alpha_s, Q^2) f_b(x_b, Q^2) \frac{d\Phi_{\text{ant}}}{d\Phi_{\text{ant}}}. 
\]

In this equation, \( \frac{d\Phi_{\text{ant}}}{d\Phi_{\text{ant}}} \) represents the antenna phase-space factorization, which provides an exact Lorentz-invariant mapping from 2 to 3 on-shell partons, that conserves global energy and momentum. Specific forms appropriate to initial–final and initial–initial antenna configurations are defined in sections [3] and [4] respectively.

The evolution variable \( Q^2 \) is a function of the phase-space point and must vanish in the unresolved limits [17]. The general formalism permits us to study different evolution variables [11, 15], though in this Letter we will restrict ourselves to a transverse-momentum type variable, defined in section [5] As in all parton showers, the description is expected to be accurate only in the strongly-ordered limit for the \( Q^2 \) of successive emissions.

The dressed or colored antenna function \( a_c \) is defined as

\[
a_c = 4\pi\alpha_s(Q^2)C\tilde{a},
\]

where \( C \) is a color factor (we recall that we use normalization conventions such that gluon and quark emission antennæ have \( C = C_A \) and \( C = 2C_F \), respectively, and gluon-splitting ones have \( C = 1 \), and \( \tilde{a} \) is a color-ordered antenna function, which embodies the factorization of QCD matrix elements in all single-unresolved soft and collinear limits. We don’t take the functions \( \tilde{a} \) to be fixed; instead we use different antenna functions with the same singular limits as one estimate of the shower uncertainty.

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\( ^1 \) Note that in [15] the normalization was \( a_c = \alpha_s/(4\pi)\tilde{a} \)
We use so-called global antenna functions \cite{4} (called sub-antenna functions with uniquely identified radiators in ref. \cite{9}) which are active over all of phase space. A backwards-evolution shower based on sector antennae in analogy to refs. \cite{18,13} is left for future work. Some, but not all, antennae needed for initial-state radiation can be chosen to be the crossings of their final–final counterparts.

An incoming particle is necessarily a hard radiator in an antenna. Therefore, the gluon emission antenna function with an incoming gluon has to reproduce the AP splitting function on its own, e. g.

\[
\bar{a}(-a,g, j,g, k_x) \frac{p_{j\rightarrow ap}}{s_{ij}} \frac{1}{1-z} P_{gg\rightarrow G}(1-z)
\]

whereas if both gluons are in the final state, the collinear singularity is reproduced by the sum of two antenna functions

\[
\bar{a}(h_x, i,g, j,g, k_x) + \bar{a}(i,g, j,g, k_x) \frac{p_{j\rightarrow ap}}{s_{ij}} \frac{1}{1-z} P_{gg\rightarrow G}(z)
\]

where the first antenna function is singular for \(i\) becoming soft, the second for \(j\).

In pure final-state showers, the \(x\) values of the incoming partons are not modified by the phase-space factorization, hence the PDF ratios in eq. (2) drop out, yielding the ordinary form of the final–final Sudakov form factor \cite{11,15}.

For initial–final antennae, only one of the PDF \(x\) values changes, and a Sudakov factor very similar to that of conventional AP showers results, with a single PDF ratio in the kernel, \(f_a(x_a, Q^2)/f_A(x_A, Q^2)\). Unlike conventional showers, however, we must also consider the backwards evolution of two initial-state partons simultaneously, generally requiring two separate parton-density factors in initial–initial antennae.

The consideration of initial–initial and initial–final antennae gives rise to one more subtlety. The basic antenna functions are color-ordered, so that in a final–final gluon-emission antenna, for example, the emitted gluon is color adjacent to both other (hard) daughter partons. That is, it is the middle parton of the color trio which is emitted. The leading-color approximation inherent in parton showers along with the symmetry of final-state phase space allows us only antennae with this ordering. When considering initial-state antennae, however, the emitted parton need not be color-adjacent to both other daughter partons; the middle parton, adjacent to both, may end up in the initial instead of the final state. We will call antennae in which the middle parton is emitted into the final state, ‘emission’ antennae; and those in which the middle parton ends up in the initial state, ‘conversion’ antennae.

For those antennae in which the type (spin) of the initial-state partons does not change after branching, we can redistribute collinear singularities to neighboring antennae so as to replace ‘conversion’ antennae by ‘emission’ antennae. For those antennae in which the type of the initial-state partons changes during branching — in which a quark backwards-evolves into a gluon or vice versa — we cannot avoid a consideration of both types of antenna function and non-emission probability.

3. Initial–Final Configurations

The pre- and post-branching partons for initial–final configurations are labeled by \(AK \rightarrow ajk\), with the other incoming parton, \(B\), acting as a passive spectator, see the illustration in fig. [1].

In general, the incoming momentum after branching will no longer be parallel to the beam direction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction. We could boost it back to the beam direction; this will transfer some of the transverse momentum generated in the emission to the rest of the event. This is the antenna analog of the recoil correction.
with $x_A/x_a = s_{AK}/(s_{AK} + s_{jk})$ and where the initial–final antenna phase space is
\[
\frac{1}{16\pi^2} \frac{s_{AK}}{(s_{AK} + s_{jk})^2} \, ds_{jk} \, ds_{aj} \tag{7}
\]
with the boundaries $0 \leq s_{jk} \leq s_{AK}(1 - x_A)/x_A$, $0 \leq s_{aj} \leq s_{AK} + s_{jk}$. We have suppressed the integration over the third coordinate of the initial–final phase space on which the emission probability does not depend.

The gluon-emission antennae can be chosen as
\[
\bar{a}(-a_q, j_g, k_q) = \frac{1}{s_{AK}} \left( \frac{2s_{ak}s_{AK}}{s_{aj} s_{jk}} + \frac{s_{jk} + s_{aj}}{s_{jk} s_{AK}} \right) \tag{8}
\]
\[
\bar{a}(-a_g, j_g, k_g) = \frac{1}{s_{AK}} \left( \frac{2s_{ak}s_{AK}}{s_{aj} s_{jk}} + \frac{s_{jk} + s_{aj}}{s_{jk} s_{AK}} \right) \tag{9}
\]
\[
\bar{a}(-a_g, j_g, k_g) = \frac{1}{s_{AK}} \left( \frac{2s_{ak}s_{AK}}{s_{aj} s_{jk}} + \frac{2s_{ak} s_{ak}}{s_{aj} (s_{ak} + s_{aj})} + \frac{2s_{jk} s_{ak}}{s_{aj} s_{AK}} + \frac{s_{aj} s_{ak}}{s_{jk} s_{AK}} \right) \tag{10}
\]
\[
\bar{a}(-a_g, j_g, b_q) = \frac{1}{s_{AK}} \left( \frac{2s_{ak}s_{AK}}{s_{aj} s_{jk}} + \frac{2s_{jk} s_{ak}}{s_{aj} (s_{ak} + s_{aj})} + \frac{2s_{jk} s_{ak}}{s_{aj} s_{AK}} + \frac{s_{aj} s_{ak}}{s_{jk} s_{AK}} \right) \tag{11}
\]
where the factor 1/2 originates from the fact that the gluon is part of two antennae.

The antenna governing the backwards-evolution of a gluon into a quark is
\[
\bar{a}(-a_q, j_g, k_q) = \frac{1}{2} \frac{s_{aj}^2 + s_{stk}^2}{s_{AK}^2} \tag{13}
\]
For the reverse process of a sea quark backwards-evolving into a gluon, we use
\[
\bar{a}_{\text{conv}}(j_q, -a_g, k_g) = \frac{1}{s_{AK}} \left( \frac{-2s_{jk} (s_{AK} - s_{aj})}{s_{aj} (s_{aj} + s_{ak})} + \frac{s_{ak}}{s_{aj}} \right) \tag{14}
\]
with a color connection $j - a - k$ at variance with the other antennae.

4. Initial–Initial Configurations

For initial–initial antennae, we label the pre- and post-branching partons by $AB \rightarrow ajb$, see fig. [1]. In the initial–initial case, we must necessarily have transverse momentum generated, which must be absorbed by the rest of the event. There are two ways of proceeding. One can allow the incoming partons to be shifted away from the beam direction after branching, and then boost back to a frame in which they are again parallel to the beam direction. Alternatively, one can fix the incoming partons to be parallel to the beam direction, and balance the new transverse momentum by boosting the rest of the event appropriately. In both cases, there is a freedom in how the longitudinal part of the emission momentum is absorbed into the initial state. This corresponds to a freedom in relating the post-branching momentum fractions $x_{a,b}$ to the pre-branching momentum fractions $x_{a,b}$. In the first case, this freedom is parametrized by the recoil or reconstruction function $r$ in combination with the Lorentz transformation boosting back to the lab frame. In the second case, it is parametrized by the functional form of $x_{a,b}$.

It turns out that these two approaches are equivalent, unlike the initial–final case. We define our recoil strategy in terms of $x_{a,b}$ here. The phase-space
factorization reads \[21\].

\[
\int \frac{dx_a}{x_a} \frac{dx_b}{x_b} d\Phi_2 (-a, -b; j, R) = \int \frac{dx_A}{x_A} \frac{dx_B}{x_B} d\Phi_1 (-A', -B'; R') d\Phi^{\mu}_{\text{ant}}
\]

(15)

with the initial–initial antenna phase space

\[
d\Phi^{\mu}_{\text{ant}} (-A', -B' \rightarrow -a, -b; j) = \frac{1}{16\pi^2 s_{ab}^2} \theta (1 - x_a) \theta (1 - x_b) ds_a ds_b
\]

(16)

where we have suppressed the integration over the angle \(\phi\) parametrizing rotations around the beam. The pre- and post-branching momenta are related by a Lorentz transform:

\[
N^{\mu}_\nu (p_R, p'_R) = g^{\mu}_\nu + \frac{2}{m_R^2} (p_R')^{\mu}(p_R)_\nu
\]

\[- \frac{2}{(p_R + p'_R)^2} (p_R + p'_R)^{\mu}(p_R + p'_R)_\nu .
\]

(17)

The phase space boundary depends directly on the definition of the post-branching momentum fractions, which is not fixed completely by the requirements of \(x_A x_B s_{AB} = x_A x_B s_{ab}\) and the behavior in the soft and collinear limits.

For gluon emission, we use

\[
x_A = \left( \frac{s_{ab} - s_{jb} s_{AB}}{s_{ab} - s_{aj} s_{ab}} \right)^{1/2},
\]

(18)

whereas for conversion, we keep one incoming momentum fixed, i.e. \(x_A/x_a = s_{AB}/s_{ab}\), \(x_b = x_B\), giving the phase space boundaries \(s_{aj} + s_{jb} \leq s_{AB}(1 - x_A)/x_A\). This corresponds to the use of a one-sided factorization, which is possible because only one collinear limit is singular, whereas the behavior of the phase space factorization in the other collinear limit is not constrained.

We use the emission antennae

\[
\tilde{a} (-a_q, j_q, -b_q) = \frac{1}{s_{AB}} \left( \frac{2 a_{ab} s_{AB}}{s_{aj} s_{jb}} + \frac{s_{aj} + s_{jb}}{s_{ab} s_{aj}} \right)
\]

(19)

\[
\tilde{a} (-a_g, j_g, -b_g) = \frac{1}{s_{AB}} \left( \frac{2 a_{ab} s_{AB}}{s_{aj} s_{jb}} + \frac{s_{aj} s_{ab}}{s_{jb} s_{ab} + s_{aj}} \right) + \frac{2 s_{aj} s_{ab}}{s_{jb} s_{ab} + s_{aj}}
\]

(20)

\[
\tilde{a} (-a_g, j_g, -b_g) = \frac{1}{s_{AB}} \left( \frac{2 a_{ab} s_{AB}}{s_{aj} s_{jb}} + \frac{s_{aj} s_{ab}}{s_{jb} s_{ab} + s_{aj}} \right) + \frac{2 s_{aj} s_{ab}}{s_{jb} s_{ab} + s_{aj}}
\]

(21)

For a quark backwards-evolving into a gluon, we use

\[
\tilde{a} (-a_g, j_g, -b_g) = \frac{1}{s_{AB}} \left( \frac{2 a_{ab} s_{AB}}{s_{aj} s_{jb}} + \frac{s_{aj} + s_{jb}}{s_{ab} s_{aj}} \right)
\]

(22)

5. Implementation and Preliminary Results

In the antenna shower, as in a conventional shower, we start the evolution at high \(Q^2\), and generate a series of branchings at successively lower \(Q^2\), stopping when we reach a shower cut-off, typically around 1 GeV. Each branching is generated according to the non-emission probability \[1\], and in this Letter we shall restrict ourselves to strict strong ordering, postponing a discussion of smooth ordering \[15\] and/or power showers \[22\] to a subsequent study.

In order to generate a branching, we must invert the function specified by the integral \[2\]. This is in general a difficult task even if the integral is double analytically, because the result involves dilogarithms. In some cases, the boundaries even make it unreasonable to perform the integral analytically. A direct inversion would in either case be quite slow. Instead, we proceed as follows. We pick a simple function — a trial antenna function \(a^{\text{trial}}\) and trial ratios of parton-density functions \(R^{\text{trial}}_{pdf}\) — which overestimates the integrand,
and veto the excess emissions generated according to the non-emission probability computed using the trial function. The trial function is chosen to capture the leading logarithmic singularities of the antenna function, and to allow the phase-space integral to be factorized into a product of one-dimensional integrals. Where possible, it is also chosen to produce an analytically invertible integral. In the final–final case, the latter requirement can always be satisfied; in initial–initial and initial–final cases, it can be satisfied for most trial antennæ. In the exceptional cases, we employ a two-stage veto. In these cases, the first-level trial function still serves to simplify the inversion of the non-emission probability computed using the trial function to maintain a reasonable efficiency [23]. Analogous issues may arise and low $Q_2$ in light-quark parton densities due to numerical instabilities. We defer their treatment to future work.

To define a concrete shower algorithm based on the above antennæ and phase-space factorizations, we have chosen to use two different evolution variables, depending on the type of antenna. For gluon emission, we use a transverse momentum, $Q^2_{\perp A} = \frac{2s_{ij}s_{jk}}{s_{ij} + s_{jk} + s_{ik}}$. For final-state branchings, $Q^2_{\perp A}$ is equal to $2p^2_{\perp A}$, the evolution measure used in Ariadne [5] (note: previous Vinçia publications used $4p^2_{\perp A}$). The maximal value of $Q^2_{\perp}$ in the final-state case is $s_{ijk}/2$. For conversion and gluon splitting antennæ, we instead use the virtuality of the only potentially singular propagator as the evolution variable. As in the final-state shower [11, 12, 15], other choices are possible within the Vinçia formalism. We defer an exploration of more general possibilities to future work.

We now turn to a few basic tests of each component of the shower algorithm. In all cases, we consider $pp$ collisions at 8 TeV CM energy, use the MSTW 2008 LO PDF set [24], with a one-loop running $\alpha_S$, normalized to $\alpha_S(m_Z) = 0.13939$. In all calculations performed here, we turn off hadronization and primordial $k_T$, as well as the underlying event, both in the Vinçia calculation and in the Pythia 8 [25] calculations to which we compare. While the evolution variables in Vinçia and Pythia are different, we have tried to match the shower cut-offs in calculations with the latter to that we use in Vinçia. This includes accounting for the difference in normalization between limiting definitions of transverse momentum. Note that while Vinçia uses a zero-mass variable flavor number scheme, Pythia 8 uses the physical quark masses everywhere. For the observables we discuss here, the effect is negligible.

In fig. 2, we show the $p_T$ spectrum of the $Z$ boson in Drell-Yan production, which is sensitive to radiation in initial–initial configurations. The main figure pane shows the peak of the distribution, while the inset shows the high-$p_T$ tail. The figure shows Vinçia curves computed using two different antenna functions: the default antenna given earlier, and an enhanced antenna function, with a finite term $-5/s_{ijk}$ added. It also shows the result obtained with Pythia 8.

The overall shape of the three curves is similar: small values at small $p_T$, rising to a peak and then declining again with a rough power-law fall in the asymptotic region, and a “knee” around $p_T \sim 90$ GeV due to the requirement of strong ordering which we have imposed here. The difference between the two Vinçia predictions illustrates the un-
Figure 2: The Drell-Yan $p_T$ spectrum. The dashed red curve shows the value computed using Vincia with default antennae functions, while the dotted green curve shows the Vincia predicted with an enhanced antenna function. The solid blue curve gives the Pythia prediction. The inset shows the high-$p_T$ tail.

Figure 3: Different color flows and corresponding emission patterns in $qq \rightarrow qq$ scattering. The straight (black) lines are quarks with arrows denoting the direction of motion in the initial or final states, and the curved (colored) lines indicating the color flow. The beam axis is horizontal, and the vertical axis is transverse to the beam. The initial-state momenta would be reversed in a Feynman diagram, so that the gluon emissions symbolically indicated by curly lines would be inside the corresponding color antennae. Forward flow is shown on the left, and backward flow on the right.

Figure 4: Angular distribution of the first gluon emission in $qq \rightarrow qq$ scattering at 45°, for the two different color flows. The light (red) histogram shows the emission density for the forward flow, and the dark (blue) histogram shows the emission density for the backward flow.
flow allows for much more radiation at $90^\circ$ than the forward one. The $p_T$ spectrum of the radiation (not shown) is also harder. The next step will be to interface the hadronization and underlying-event models in Pythia, and compare to experimental studies, such as the one by CDF \cite{27} (we note that an update of that study, correcting it to the hadron level, would be highly useful to the MC community).

Finally, to demonstrate the combination of all shower components acting together, we show the dijet decorrelation angle (the azimuthal angle between the two leading jets), $\Delta \phi_{JJ}$, in fig. \ref{fig:decorrelation}. Using FastJet \cite{28}, we consider anti-$k_T$ jets with radius parameter $R = 0.4$. We demand the two leading jets to have transverse momentum above 100 GeV and to be at rapidities $|y| < 2.8$. Note that Pythia 8 produces about 40% more jet events which pass these cuts than Vincia both with or without enhanced antenna functions, partly due to its more active recoil strategy, which allows the original dijet system to build up transverse momentum successively during the shower cascade. This difference is not visible in fig. \ref{fig:decorrelation} as the distributions are all normalized to unity. The two Vincia distributions are broadly similar to the Pythia 8 distribution, and very similar in the two-jet region ($\Delta \phi_{JJ} \sim \pi$) where the parton-shower approximation should be reliable. The differences are substantial in the region below $\Delta \phi_{JJ} < 3\pi/4$, where hard real emission is important. In this region, fixed-order calculations should be reliable, but unmatched parton showers will not be. Nonetheless, the difference between the Vincia calculation with default antenna strength and the Pythia 8 calculation is similar to that between the two Vincia calculations, suggesting again that the variation provides a good qualitative assessment of the uncertainty. We defer a comparison of this distribution with fixed-order calculations to future studies.

6. Conclusion and Outlook

In this Letter, we have presented the outline of a formalism for an antenna shower for hadron collisions, along with results from an initial implementation as a plug-in to Pythia 8. The formalism requires the introduction of new antennae, corresponding to one or both parents being initial-state partons. These should be further subdivided into the new categories of emission and conversion antennae based on their color flows. The formalism also requires factorizations suitable for phase spaces involving initial-state partons, and introduces ratios of parton densities into the non-emission probabilities governing the shower evolution.

We have chosen to implement the shower as a plug-in to the Pythia 8 program, which takes advantage of the latter’s flexible framework and makes use of its utilities, structures, and overall management of the branching process. In this approach, it replaces Pythia’s shower with an antenna shower. For practical and efficiency reasons, we uniformly adopt a trial-and-veto algorithm for generating branchings. The trial functions used in the implementation will be described elsewhere.

We expect to implement further optimizations of the branching step in future work. The leading-order matching approach described in ref. \cite{15} should carry over to the initial-state showering described here, and will be an important next step for the development of Vincia.
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