Phenomenological study of the interplay between IR-improved DGLAP-CS theory and the precision of an NLO ME matched parton shower MC

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HIGHLIGHTS

- Using LHC data, we show that IR-improved DGLAP-CS kernels with exact NLO Shower/ME matching improves MC precision.
- We discuss other possible approaches in comparison with ours.
- We propose experimental tests to discriminate between competing approaches.

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ABSTRACT

We present a phenomenological study of the current status of the application of our approach of *exact* amplitude-based resummation in quantum field theory to precision QCD calculations, by realistic MC event generator methods, as needed for precision LHC physics. We discuss recent results as they relate to the interplay of the attendant IR-improved DGLAP-CS theory of one of us and the precision of exact NLO matrix-element matched parton shower MC's in the Herwig6.5 environment as determined by comparison to recent LHC experimental observations on single heavy gauge boson production and decay. The level of agreement between the new theory and the data continues to be a reason for optimism. In the spirit of completeness, we discuss as well other approaches

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to the same theoretical predictions that we make here from the standpoint of physical precision with an eye toward the (sub-)1% QCD ⊗ EW total theoretical precision regime for LHC physics.

1. Introduction

With the recent announcement [1] of a Brout–Englert–Higgs (BEH) [2] candidate boson after the start-up and successful running of the LHC for three years, we have entered the era of precision QCD, by which we mean predictions for QCD processes at the total precision tag of 1% or better. Given the expected role of precision comparison between theory and experiment in determining the detailed properties of the newly discovered BEH boson candidate, the attendant need for exact, amplitude-based resummation of large higher order effects is now more paramount. We have argued in Refs. [3,4] that such resummation allows one to have better than 1% theoretical precision as a realistic goal in such comparisons, so that one can indeed distinguish new physics (NP) from higher order SM processes and can distinguish different models of new physics from one another as well. In what follows, with an eye toward its interplay with NLO exact, matrix element matched parton shower MC precision issues, we present the status of this approach to precision QCD for the LHC in connection with its attendant IR-improved DGLAP-CS [5,6] theory [7,8] realization via HERWIRI1.031 [9] in the HERWIG6.5 [10] environment. We realize here the attendant exact, NLO matrix element matched parton shower MC realizations for both HERWIRI1.031 and HERWIG6.5 in the MC@NLO [11] methodology, wherewith we make the corresponding comparisons with recent LHC data that we present herein.

In this way, we continue the strategy of building on existing platforms to develop and realize a path toward precision QCD for the physics of the LHC. What we exhibit is an explicit union of the new IR-improved DGLAP-CS theory and the MC@NLO realization of exact NLO matrix element (ME) matched parton shower MC theory. We note that we are also pursuing the implementation [12] of the new IR-improved DGLAP-CS theory for HERWIG++ [13], HERWIRI++, for PYTHIA8 [14] and for SHERPA [15], as well as the corresponding NLO ME/parton shower matching realizations in the POWHEG [16] framework. For, our ultimate goal is a provable precision tag on our theoretical predictions and one of the strongest cross checks on theoretical precision is the difference between two independent realizations of the attendant theoretical calculation. We will present such cross checks elsewhere [12].

The paper is organized as follows. We motivate the development of the new IR-improved DGLAP-CS theory in our approach to precision QCD theory by showing in the next section how it arises naturally in the effort to realize a provable precision from our approach [4] to precision LHC physics. This will allow us to expose properly the interplay between the NLO ME matched parton shower MC precision and the new IR-improved DGLAP-CS theory. In the interest of completeness, we also review in the next section the relevant aspects of the approach in Ref. [4], which is an amplitude-based QED ⊗ QCD(≡QCD ⊗ QED)[1] exact resummation theory realized by MC methods. This paper is an extension of a shorter report of our results and analysis which is given in Ref. [17]. Accordingly, also in the interest of completeness, we discuss as well in the next section other possible approaches to the predictions we make herein from the standpoint of their physical precision. We believe this discussion will aid the reader in putting the results we present in a proper context. Section 3 contains the applications to the recent LHC data on single heavy gauge boson production with an eye toward the analysis in Ref. [9] of the analogous processes at the Tevatron. In this paper we will focus on the single

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[1] In our resummation algebra, only those terms that commute in the action of QED and QCD interactions are resummed in the attendant resummation exponents (see Eq. (4)) so that we get the same results order-by-order in α and α_s independent of which set of large infrared corrections are resummed first.

$Z/\gamma*$ production and decay to lepton pairs for definiteness. The other heavy gauge boson processes will be taken up elsewhere [12]. Section 4 contains our concluding observations and our outlook.

2. Review of our approach to precision LHC physics

We take the starting point for what we discuss here as the following fully differential representation of a hard LHC scattering process:

$$d\sigma = \sum_{i,j} \int dx_1 dx_2 F_i(x_1) F_j(x_2) d\hat{\sigma}_{\text{res}}(x_1 x_2 s), \qquad (1)$$

where the $\{F_j\}$ and $d\hat{\sigma}_{\text{res}}$ are the respective parton densities and reduced hard differential cross section and where we use the subscript on the latter to indicate that it has been resummed for all large EW and QCD higher order corrections in a manner consistent with achieving a total precision tag of 1% or better for the total theoretical precision of (1). The key theoretical issue for precision QCD for the LHC is then the proof of the correctness of the value of the total theoretical precision $\Delta\sigma_{\text{th}}$ of (1). The theoretical precision in question can be represented as follows:

$$\Delta\sigma_{\text{th}} = \Delta F \oplus \Delta\hat{\sigma}_{\text{res}} \qquad (2)$$

where ΔA is the contribution of the uncertainty on the quantity A to $\Delta\sigma_{\text{th}}$. Here, we stress that we discuss the situation in which the two errors in (2) are independent for definiteness; (2) has to be modified accordingly when they are not. We further stress that the proof of the correctness of the value of the total theoretical precision $\Delta\sigma_{\text{th}}$ is indispensable in order to validate the application of a given theoretical prediction to precision experimental observations, for the discussion of the signals and the backgrounds for both Standard Model (SM) and new physics (NP) studies, and more specifically for the overall normalization of the cross sections in such studies. If one uses a calculation with an unknown value of $\Delta\sigma_{\text{th}}$ for the attendant studies, the NP itself can be missed. We do feel a need to continue to state here that this point simply cannot be emphasized too much.

In the interest of completeness, we note that, by our definition, which follows the discussion in Ref. [18], $\Delta\sigma_{\text{th}}$ is the total theoretical uncertainty that derives from the physical and technical precision contributions: the physical precision contribution, $\Delta\sigma_{\text{th}}^{\text{phys}}$, arises from such sources as missing graphs, approximations to graphs, truncations, …; the technical precision contribution, $\Delta\sigma_{\text{th}}^{\text{tech}}$, arises from such sources as bugs in codes, numerical rounding errors, convergence issues, etc. We want to observe that here, when we reference bugs in codes we have in mind that all gross errors such as those that give obviously wrong results, as determined by cross checks, are eliminated and we have left programming errors such as those in the logic: suppose for programming error reasons a DO-loop ends at 999 steps instead of the intended 1000 steps, resulting in a per mille level error, that could alternate in sign from event to event. As per mille level accuracy is good enough in many applications, the program would remain reliable, but it would have what we call a technical precision error at the per mille level. With this understanding, the total theoretical error is then given by

$$\Delta\sigma_{\text{th}} = \Delta\sigma_{\text{th}}^{\text{phys}} \oplus \Delta\sigma_{\text{th}}^{\text{tech}}. \qquad (3)$$

Although the desired value for $\Delta\sigma_{\text{th}}$ depends on the specific requirements of the observations, as a general rule, it should fulfill $\Delta\sigma_{\text{th}} \leq f \Delta\sigma_{\text{expt}}$, where $\Delta\sigma_{\text{expt}}$ is the respective experimental error and $f \lesssim \frac{1}{2}$, so that the theoretical uncertainty does not significantly adversely affect the attendant physics analysis of the data.

In the interest of completeness, we will now recapitulate the theory we have developed to realize such precision rigorously. This theory is the QCD $\otimes$ QED resummation theory in Refs. [4] for the reduced cross section in (1) and for the resummation of the evolution of the parton densities therein as well. More precisely, the master formula, which applies both to the resummation of the reduced

cross section and to that of the evolution of the parton densities, may be identified as

$$
\begin{aligned}
d\bar{\sigma}_{\text{res}} = e^{\text{SUM}_{\text{IR}}(\text{QCED})} \sum_{n,m=0}^{\infty} \frac{1}{n!m!} \int \prod_{j_1=1}^{n} \frac{d^3 k_{j_1}}{k_{j_1}} \\
\times \prod_{j_2=1}^{m} \frac{d^3 k'_{j_2}}{k'_{j_2}} \int \frac{d^4 y}{(2\pi)^4} e^{iy\cdot(p_1+q_1-p_2-q_2-\sum k_{j_1}-\sum k'_{j_2})+D_{\text{QCED}}} \\
\times \tilde{\bar{\beta}}_{n,m}(k_1,\ldots,k_n; k'_1,\ldots,k'_m) \frac{d^3 p_2}{p_2^0} \frac{d^3 q_2}{q_2^0},
\end{aligned}
\tag{4}
$$

where $d\bar{\sigma}_{\text{res}}$ is either the reduced cross section $d\hat{\sigma}_{\text{res}}$ or the differential rate associated to a DGLAP-CS [5,6] kernel involved in the evolution of the $\{F_j\}$ and where the *new* (YFS-style [19,20]) *non-Abelian* residuals $\tilde{\bar{\beta}}_{n,m}(k_1,\ldots,k_n; k'_1,\ldots,k'_m)$ have n hard gluons and m hard photons and we show the final state with two hard final partons with momenta p_2, q_2 specified for a generic 2f final state for definiteness. The infrared functions $\text{SUM}_{\text{IR}}(\text{QCED})$, D_{QCED} are defined in Refs. [4,7,8] as follows:

$$
\text{SUM}_{\text{IR}}(\text{QCED}) = 2\alpha_s \Re B_{\text{QCED}}^{nls} + 2\alpha_s \tilde{B}_{\text{QCED}}^{nls}
$$

$$
D_{\text{QCED}} = \int \frac{d^3 k}{k^0} \left(e^{-iky} - \theta(K_{\max} - k^0)\right) \tilde{S}_{\text{QCED}}^{nls}
\tag{5}
$$

where the dummy parameter $K_{\max}$ is such that nothing depends on it and where we have introduced

$$
\begin{aligned}
B_{\text{QCED}}^{nls} &\equiv B_{\text{QCD}}^{nls} + B_{\text{QED}}^{nls}, \\
\tilde{B}_{\text{QCED}}^{nls} &\equiv \tilde{B}_{\text{QCD}}^{nls} + \tilde{B}_{\text{QED}}^{nls}, \\
\tilde{S}_{\text{QCED}}^{nls} &\equiv \tilde{S}_{\text{QCD}}^{nls} + \tilde{S}_{\text{QED}}^{nls}.
\end{aligned}
\tag{6}
$$

Here, the superscript *nls* denotes that the infrared functions are DGLAP-CS synthesized as explained in Refs. [4,7,8,21] and the infrared functions B_A, $\tilde{B}_A$, $\tilde{S}_A$, $A = $ QCD, QED, are given in Refs. [4,7,8,19,20]. We stress that simultaneous resummation of QED and QCD large IR effects is exact here.

We may describe the physical meanings of the key components in the master formula (4) as follows. In the language of Ref. [22], the exponent $\text{SUM}_{\text{IR}}(\text{QCED})$ sums up to the infinite order the maximal leading IR singular terms in the cross section for soft emission below a dummy parameter $K_{\max}$ and the exponent D_{QCED} does the same for the regime above $K_{\max}$ so that (4) is independent of $K_{\max}$—it cancels between $\text{SUM}_{\text{IR}}(\text{QCED})$ and D_{QCED}.[2] In order to maintain exactness order by order in perturbation theory in both α and α_s in the presence of these resummed terms, as explained in Refs. [4,7,8] we generate the residuals $\tilde{\bar{\beta}}_{n,m}$ by iterative computation to match the attendant exact results to all orders in α and α_s. We need to emphasize the following most important point. In our formulation in (4) *the entire soft gluon phase space is included in the representation—no part of it is dropped.*

As it is explained in Refs. [4], the new non-Abelian residuals $\tilde{\bar{\beta}}_{m,n}$ allow rigorous shower/ME matching via their shower subtracted analogs: in (4) we make the replacements

$$
\tilde{\bar{\beta}}_{n,m} \rightarrow \hat{\tilde{\bar{\beta}}}_{n,m}
\tag{7}
$$

where the $\hat{\tilde{\bar{\beta}}}_{n,m}$ have had all effects in the showers associated to the $\{F_j\}$ removed from them. To see how we make contact between the $\hat{\tilde{\bar{\beta}}}_{n,m}$ and the differential distributions in MC@NLO we proceed as

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[2] If we want to include more of the maximal exponentiating terms from the formalism of Ref. [22] in the two exponents $\text{SUM}_{\text{IR}}(\text{QCED})$, D_{QCED}, we may do so with a consequent change in the attendant residuals $\tilde{\bar{\beta}}_{n,m}$.

follows. We represent the MC@NLO differential cross section via [11]

$$d\sigma_{\text{MC@NLO}} = \left[B + V + \int (R_{\text{MC}} - C) d\Phi_R \right] d\Phi_B \left[\Delta_{\text{MC}}(0) + \int (R_{\text{MC}}/B) \Delta_{\text{MC}}(k_T) d\Phi_R \right]$$
$$+ (R - R_{\text{MC}}) \Delta_{\text{MC}}(k_T) d\Phi_B d\Phi_R \tag{8}$$

where B is Born distribution, V is the regularized virtual contribution, C is the corresponding counter-term required at exact NLO, R is the respective exact real emission distribution for exact NLO, $R_{\text{MC}} = R_{\text{MC}}(P_{AB})$ is the parton shower real emission distribution so that the Sudakov form factor is

$$\Delta_{\text{MC}}(p_T) = e^{\left[- \int d\Phi_R \frac{R_{\text{MC}}(\Phi_B, \Phi_R)}{B} \theta(k_T(\Phi_B, \Phi_R) - p_T) \right]},$$

where as usual it describes the respective no-emission probability. The respective Born and real emission differential phase spaces are denoted by $d\Phi_A, A = B, R$, respectively. We may note further that the representation of the differential distribution for MC@NLO in (8) is an explicit realization of the compensation between real and virtual divergent soft effects discussed in the Appendices of Refs. [7,8] in establishing the validity of (4) for QCD—all of the terms on the RHS of (8) are infrared finite. Indeed, from comparison with (4) restricted to its QCD aspect we get the identifications, accurate to $\mathcal{O}(\alpha_s)$,

$$\frac{1}{2} \hat{\bar{\beta}}_{0,0} = \bar{B} + (\bar{B}/\Delta_{\text{MC}}(0)) \int (R_{\text{MC}}/B) \Delta_{\text{MC}}(k_T) d\Phi_R$$
$$\frac{1}{2} \hat{\bar{\beta}}_{1,0} = R - R_{\text{MC}} - B\tilde{S}_{\text{QCD}} \tag{9}$$

where we defined [11]

$$\bar{B} = B(1 - 2\alpha_s \Re B_{\text{QCD}}) + V + \int (R_{\text{MC}} - C) d\Phi_R$$

and we understand here that the DGLAP-CS kernels in R_{MC} are to be taken as the IR-improved ones as we exhibit below [7,8]. Here we have written the QCD virtual and real infrared functions B_{QCD} and $\tilde{S}_{\text{QCD}}$ respectively without the superscript *nls* for simplicity of notation and they are understood to be DGLAP-CS synthesized as explained in Refs. [4,7,8] so that we avoid double counting of effects. We also re-emphasize that we do not drop any effects here in (9). We observe further that, in view of (9), the way to the extension of frameworks such as MC@NLO to exact higher orders in $\{\alpha_s, \alpha\}$ is therefore open via our $\hat{\bar{\beta}}_{n,m}$ and will be taken up elsewhere [12].

We point out that in Refs. [7–9] our methods for resummation of the QCD theory have been shown to be fully consistent with the methods in Refs. [23,24]. Specifically, it is shown in Refs. [7–9] that the methods in Refs. [23,24] give approximations to our hard gluon residuals $\hat{\bar{\beta}}_{n,0}$; for unlike the master formula in (4), the methods in Refs. [23,24] are not exact results. To see this, observe that the threshold-resummation methods in Refs. [23], using the result that, for any function $f(z)$,

$$\left| \int_0^1 dz z^{n-1} f(z) \right| \leq \left(\frac{1}{n} \right) \max_{z \in [0,1]} |f(z)|,$$

drop non-singular contributions to the cross section at $z \to 1$ in resumming the logs in n-Mellin space. The SCET theory in Refs. [24] drops terms of $\mathcal{O}(\lambda)$ at the level of the amplitude, where $\lambda = \sqrt{\Lambda/Q}$ for a process with the hard scale Q with $\Lambda \sim 0.3$ GeV so that, for $Q \sim 100$ GeV, we have $\lambda \cong 5.5\%$. From the known equivalence of the two approaches, the errors in the threshold resummation must be similar. Evidently, we can only use these approaches as a guide to our new non-Abelian residuals as we develop results for the (sub-)1% precision regime.

In view of the specific processes which we consider in the next section, we will continue here with our consistency discussion as it relates to the theory of QCD resummation for the specific heavy gauge boson production-type processes offered in Refs. [25,26], which is again an approximate formalism

that could be used to make approximations to our hard gluon residuals $\hat{\tilde{\beta}}_{n,0}$. The theory in Refs. [25,26] is in wide use at the LHC and in the data analyses for the Tevatron—see for example the recent analyses in Refs. [27] where this theory, as it is implemented in the MC integration program RESBOS [28–30] is compared to recent LHC data and to recent analyses of Tevatron data. We need to note that the theory in Refs. [25,26] builds on and extends beyond the considerable literature in Refs. [31] aimed at the analogous processes to those under study here. Could we perhaps employ this formalism in Refs. [25,26] as well to reach ≤1% physical precision QCD predictions? Let us note that the authors in Ref. [26] give us a hint to the answer to our question in their footnote on the journal page with number 215 for that paper, wherein they equate as 'negligible' 20% in discussing possible nonperturbative contributions in their formalism. Let us keep this footnote in mind in what follows.

We first note that a defining formula for the approach in Refs. [25,26] is that for the differential cross section for the p_T distribution for the heavy gauge boson production in hadron–hadron collisions, which we specialize here to the case of the Drell–Yan $\gamma*$ production for definiteness,[3] where we record the result in Eq. (2.2) from Ref. [26]:

$$\frac{d\sigma}{dQ^2 dy dQ_T^2} \sim \frac{4\pi^2\alpha^2}{9Q^2 s}\left\{\int \frac{d^2 b}{(2\pi)^2} e^{i\vec{Q}_T \cdot \vec{b}} \widetilde{W}(b; Q, x_A, x_B) + Y(Q_T; Q, x_A, x_B)\right\} \qquad (10)$$

where we have the usual kinematics so that $\vec{Q}_T = \vec{p}_T$ is the $\gamma*$ transverse momentum, A, B are protons at the LHC, s is the cms squared energy of the protons, Q^μ is the $\gamma*$ 4-momentum so that Q^2 is the $\gamma*$ mass squared, and $y = \frac{1}{2}\ln(Q^+/Q^-)$ is the $\gamma*$ rapidity so that $x_A = e^y Q/\sqrt{s}$ and $x_B = e^{-y}Q/\sqrt{s}$. We have in mind that Q is near M_Z here. In (10), the term involving $\widetilde{W}$ carries the effects from QCD resummation as developed in Refs. [25,26] and the Y term includes those contributions which are 'regular' at $Q_T = p_T \to 0$ in the sense of Refs. [25,26], i.e., order by order in perturbation theory they are derived from the parts of the attendant hard scattering coefficients that are less singular than $Q_T^{-2} \times$ (logs or 1) or $\delta(\vec{Q}_T)$ as $Q_T = p_T \to 0$. Our question concerns the physical precision of the $\widetilde{W}$ term; for the Y term is perturbative and can be computed in principle to the required accuracy by the standard methods.

The result for $\widetilde{W}$ given in Refs. [25,26] is as follows. When we have $b \ll 1/\Lambda$ where Λ is a typical QCD hadronic mass scale such as the inverse of the proton radius, the result for $\widetilde{W}$ is [25,26]

$$\widetilde{W}(b; Q, x_A, x_B) = \sum_{a,b} \int_{x_A}^1 d\xi_A \int_{x_B}^1 d\xi_B F_{a/A}(\xi_A; \mu) F_{b/B}(\xi_B; \mu) e^{-S_{sc}}$$
$$\times \sum_j e_j^2 C_{ja}(x_A/\xi_A, b; C_1/C_2; g(\mu), \mu) C_{jb}(x_B/\xi_B, b; C_1/C_2; g(\mu), \mu), \qquad (11)$$

where we define [25,26]

$$S_{sc} = \int_{C_1^2/b^2}^{C_2^2 Q^2} \frac{d\bar{\mu}^2}{\bar{\mu}^2}\left[\ln\left(\frac{C_2^2 Q^2}{\bar{\mu}^2}\right) A(g(\bar{\mu}); C_1) + B(g(\bar{\mu}); C_1, C_2)\right]. \qquad (12)$$

Here we show explicitly the dependences of the parton density functions $\{F_\ell\}$ (note that our $F_\ell(x)$ corresponds to $f_\ell(x)/x$ in Refs. [25,26]), and we show as well the dependences of the perturbatively calculable exponentiation and scattering coefficient functions A, B, and C as they are given in Refs. [25,26]. The scale μ is usually set to C_1/b and C_1 and C_2 are order 1 constants chosen to optimize the resultant perturbation expansions—see Refs. [25,26] for more discussion on this point. The authors in Refs. [25,26] note that this result for $\widetilde{W}(b)$, when $1/Q \ll b \ll 1/\Lambda$, is accurate up to terms $\mathcal{O}(m_q/Q, 1/(bQ))$ where m_q represents the quark masses. In addition, the latter authors argue that in

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[3] The analogous results for the $W^\pm$ and $Z/\gamma*$ are obtained by straightforward substitutional manipulations of the EW aspects of the formula we record here as described in Ref. [26] so that we omit such manipulations here without loss of content of the QCD aspects our discussion.

the result (10) they have dropped terms of $\mathcal{O}(m_q/Q, Q_T/Q)$ in the regime where $0 \le Q_T \ll Q$. These last two statements are seen explicitly in Sect. 9 of Ref. [25] in Eq. (9.1) and the equation immediately preceding it, for reference.

When we have to consider as well the regime $b \gtrsim 1/\Lambda$, the authors in Refs. [25,26] argue that we have to replace the functions A, B, and C according to

$$A(g(\bar{\mu}); C_1) \to A(g(\bar{\mu}), m_q/\bar{\mu}; C_1)$$

$$B(g(\bar{\mu}); C_1, C_2) \to B(g(\bar{\mu}), m_q/\bar{\mu}; C_1, C_2)$$

$$\sum_{a,b} \int_{x_A}^1 d\xi_A \int_{x_B}^1 d\xi_B F_{a/A}(\xi_A; \mu) F_{b/B}(\xi_B; \mu) \sum_j e_j^2 C_{ja}(x_A/\xi_A, b; C_1/C_2; g(\mu), \mu)$$

$$\times C_{jb}(x_B/\xi_B, b; C_1/C_2; g(\mu), \mu) \to \sum_j e_j^2 \bar{\mathcal{P}}_{j/A}(x_A, b; C_1/C_2) \bar{\mathcal{P}}_{j/B}(x_B, b; C_1/C_2) \quad (13)$$

where we see that we can no longer neglect the quark masses and that we no longer have the convolutions $C * F$: the functions $\bar{\mathcal{P}}_{j/A}$ reduce to the latter form when we take $b \ll 1/\Lambda$ and drop terms $\mathcal{O}(b\Lambda)$, according to the arguments in Refs. [25,26].

The authors in Refs. [25,26] then join these two results, that for the regime $b \ll 1/\Lambda$ and that for the regime $b \gtrsim 1/\Lambda$, via an ansatz as follows. Defining

$$\widetilde{W}(b; Q, x_A, x_B) \equiv \sum_j e_j^2 \widetilde{W}_j(b; Q, x_A, x_B), \quad (14)$$

which we see from (11) and (13) is well-defined for both regimes of interest in b, the authors in Ref. [26] use

$$b^* = b/\sqrt{1 + b^2/b_{\max}^2} \quad (15)$$

where $b_{\max}$ is a parameter still in the perturbative regime where (11) holds, to write the joining ansatz

$$\widetilde{W}_j(b; Q, x_A, x_B) = \widetilde{W}_j(b^*; Q, x_A, x_B) e^{\{-\ln(Q^2/Q_0^2)g_1(b) - g_{j/A}(x_A, b) - g_{j/B}(x_B, b)\}}, \quad (16)$$

where the functions $g_1, g_{j/A}, g_{j/B}$ have non-perturbative content and must be determined from data. With this ansatz, the authors in Ref. [26] arrive at the representation

$$\frac{d\sigma}{dQ^2 dy dQ_T^2} \sim \frac{4\pi^2 \alpha^2}{9Q^2 s} \left\{ \int \frac{d^2 b}{(2\pi)^2} e^{i\vec{Q}_T \cdot \vec{b}} \sum_j e_j^2 \widetilde{W}_j(b^*; Q, x_A, x_B) \right.$$

$$\left. \times e^{\{-\ln(Q^2/Q_0^2)g_1(b) - g_{j/A}(x_A, b) - g_{j/B}(x_B, b)\}} + Y(Q_T; Q, x_A, x_B) \right\}. \quad (17)$$

We point-out for completeness and illustration that in RESBOS in Refs. [29,30] three realizations of the non-perturbative functions in (17) were considered as follows:

$$\ln(Q^2/Q_0^2)g_1(b) + g_{j/A}(x_A, b) + g_{j/B}(x_B, b)$$

$$= \begin{cases} [g_1 + g_2 \ln(Q/(2Q_0))]b^2, & \text{DWS [32]} \\ [g_1 + g_2 \ln(Q/(2Q_0))]b^2 + g_1 g_3 \ln(100 x_A x_B)b, & \text{LY [29]} \\ [g_1 + g_2 \ln(Q/(2Q_0)) + g_1 g_3 \ln(100 x_A x_B)]b^2, & \text{BLNY [30]}. \end{cases} \quad (18)$$

We observe that the g_i are parameters on the RHS of (18) and that the best fit to the data considered in Ref. [30] was obtained from the BLNY parametrization with the values

$$g_1 = 0.21^{+0.01}_{-0.01} \text{ GeV}^2, \qquad g_2 = 0.68^{+0.01}_{-0.02} \text{ GeV}^2, \qquad g_3 = -0.6^{+0.05}_{-0.04} \quad (19)$$

when $b_{\max} = 0.5 \text{ GeV}^{-1}$ and $Q_0 = 1.6 \text{ GeV}$.

On the question of the physical precision of (17), we first observe that, with the errors shown on the constants g_i for the best BLNY parametrization, the -0.02 error on g_2 represents already a 1.5% uncertainty at the saddle point position for the integration over b in the respective Z production analog of (10) found in Ref. [26] and recapitulated below in (20)—this alone exceeds the theory error budget in precision QCD theory that we advocate here so that one would need these errors reduced considerably for them to be useful for our purposes. But, what is more important is that, on the LHS of (18) the g_ℓ which do not (do) multiply $\ln(Q/Q_0)$ are unspecified *functions* of x_j, b (b) that are required to vanish at $b = 0$ while on the RHS these dependences are simplified to second order polynomials in b and either no or linear $\ln(x_A x_B)$ dependence on x_j: these simplifications are generically ad hoc and cannot be considered as a rigorous platform for testing the fundamental QCD theory. They (the simplifications) do not follow from the formalism of Refs. [25,26]. As the authors in Ref. [32] have emphasized, the latter formalism could only be used to give a prediction for perturbative QCD in the regime where the details of the non-perturbative parametrization are not important. We cannot emphasize this point too much.[4]

On the question of the physical precision of (17) we additionally note that in the regime $0 \leq Q_T \ll Q$ it has an error of $\mathcal{O}(Q_T/Q)$ according to the authors in Refs. [25,26], as such terms have been dropped, order by order in perturbation theory in the construction of $\tilde{W}_j$—we make reference again specifically to Eq. (9.1) and the equation immediately preceding it in Ref. [25]. We stress here that in any given realization of (17) such as that in Ref. [28] the quantity Y is only known to a finite order in α_s, in the latter reference this is to first order in α_s. This means that the physical precision error $\mathcal{O}(Q_T/Q)$ on the part of the resummed singular part of (17) from $\mathcal{O}(\alpha_s^n)$, $n \geq 2$, cannot be compensated by Y in Ref. [28], for example. If Y is known to $\mathcal{O}(\alpha_s^n)$, physical precision error $\mathcal{O}(Q_T/Q)$ from the $\tilde{W}_j$ part of (17) from $\mathcal{O}(\alpha_s^m)$, $m \geq n + 1$, cannot be compensated by Y even in principle. For $Q = M_Z$ and $Q_T = 5$ GeV, this is an error $\cong 5.5\%$ and its twice this size at $Q_T = 10$ GeV. This shows that the formalism in Refs. [25,26] cannot be used for (sub-)1% theory predictions for the heavy gauge boson p_T spectrum at the LHC. One might think that as $Q_T \to 0$, where these dropped terms should vanish, the theory in Refs. [25,26] would be accurate enough for (sub-)1% precision QCD predictions. This is not true as well because at $Q_T = 0$, following Parisi and Petronzio in Refs. [31], the authors in Ref. [26] show that the integral over $\vec{b}$ in (10) is dominated by a saddle point at

$$b_{SP} = \frac{1}{\Lambda}\left(\frac{Q}{\Lambda}\right)^{-A^{(1)}/(A^{(1)}+\beta_1)} \tag{20}$$

where we have $A = \sum_N A^{(N)}(\alpha_s/\pi)^N$ and $\beta_1 = \frac{1}{12}(33 - 2N_F)$ where N_F is the number of effective quark flavors, which we take as 5 here. Thus, we have [25,26] $A^{(1)} = 4/3$ so that

$$b_{SP} = \frac{1}{\Lambda}\left(\frac{Q}{\Lambda}\right)^{-0.41}.$$

For $Q = M_Z$, we get $b_{SP} \cong 0.48$ GeV^{-1} when we follow Ref. [26] and take $\Lambda = 0.15$ GeV. This result for b_{SP} is in the perturbative regime. Moreover, we know from Refs. [25,26] that the error on the Z production analog of $\tilde{W}_j$ at this point is $\mathcal{O}(1/(b_{SP}M_Z)) \cong 2.3\%$—see again Eq. (9.1) in Ref. [25]. Thus, even when the perturbative regime obtains so that the result (11) should be reliable, the error in it is too large for use as anything but a guide in constructing our precision QCD $\otimes$ EW theory residuals $\tilde{\tilde{\beta}}_{n,m}$ in our master formula (4) for $\lesssim$1% precision LHC physics. The results reported in Refs. [27] on the

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[4] Note that, with the IR-improvement of the parton shower in Herwiri1.031, there is no longer any need even for an IR cut-off parameter, as all formerly IR-divergent integrals now converge; there is no issue of the validity of perturbation theory in the shower because the shower resums the large logs in it to all orders in α_s; the shower runs down to the hadronization scale ∼1 GeV and the transition to hadrons is a smooth one due to local parton hadron duality [33] and the precociousness of Bjorken scaling [34,35] which occurs already at 1_+ GeV^2—the properties of the shower such as its p_T and its energy are independent of the transition to hadrons to very high accuracy. No arbitrary ad hoc parameters that affect resummed results from the shower such as the parameters in (18) appear—even the un-IR-improved parton shower in Herwig 6.5 is well-behaved at small p_T and no part of the result of $\mathcal{O}(Q_T/Q)$ in (4) is just dropped as it is in (17).

comparison of RESBOS with recent ATLAS and Tevatron data related to the p_T spectrum in single Z/γ^* are consistent with our estimates on the physical precision of (17): RESBOS misses the data by 2% for the p_T near 0 and it misses the data by $\gtrsim$5% for the regime of $p_T \gtrsim$ 10 GeV.

The non-perturbative factor in (17) compromises the predictive power of the formalism in Refs. [25,26]. The authors in Ref. [32] have argued that for W, Z production as considered here the prediction in (17) in the regime 6 GeV $< Q_T <$ 16 GeV is insensitive to the non-perturbative parametrization in (18). This leaves the effects we have discussed above as the main obstacles to using the theory in Refs. [25,26] for precision QCD physics predictions at the LHC in this latter regime when that precision is at or below 1%.

We note that in Ref. [36], another version of (10) is presented which has the same physical precision limits as that realized in RESBOS with a different treatment of the non-perturbative regime. As the comparison with the data shown in Refs. [27] shows in the perturbative regime, our estimates of this physical precision for this regime also apply to the realization of (10) in Ref. [36]: it cannot be used for 1% precision LHC physics studies.

Here, let us note that the authors in Ref. [36] discuss the uncertainty in their results as a function of the variation of the (perturbative scale) values associated with their renormalization, factorization and resummation scales. For example, they estimate that this is at the level of 10% near the peak of the Q_T spectrum in the single $Z/\gamma*$ production at the LHC. These uncertainties can in principle be reduced by computing the perturbative terms in their results to higher and higher orders. They also estimate an additional PDF error at this region at the level of 2%, which in principle can be reduced by improving the determination of the respective PDF's. We stress that the error in the defining result (10) derived in Refs. [25,26] on which their results are based that we discuss here is an error of $\mathcal{O}(Q_T/Q)$ that applies order by order in the perturbation theory—it is separate from the scale and PDF errors discussed in Ref. [36].

Indeed, in Ref. [37], the SCET approach is used to recover (10) and it is shown explicitly that the improvement of the perturbative scale errors indeed occurs when higher order corrections are included in the calculation of the perturbative terms in the respective SCET realization of (10). We stress again that the error that we discuss here of $\mathcal{O}(Q_T/Q)$ due to the approximations in the defining derivation of (10) in Refs. [25,26] also applies to the order-by-order results in Refs. [37]. As we noted above, SCET involves for single Z production at the LHC the defining error $\sqrt{\Lambda/M_Z} \sim$ 5.7% for the typical hadronic transverse size $\Lambda \cong$ 0.3 GeV and this is consistent with the approximations made in Refs. [25,26]. Such approximations cannot be used for the 1% precision QCD theory that we have as our goal for (4) here.

In summary, this last remark is beginning to be manifest as we see in Refs. [27] in the comparison between the recent Tevatron and LHC data on the Q_T spectra in the $Z/\gamma*$ production and the predictions of RESBOS and of Ref. [36]. Indeed, even though a new Q_T-related variable [38] is used in some of the comparisons, $\phi_\eta^* = \tan(\frac{1}{2}(\pi - \Delta\phi)) \sin\theta^* \cong \left| \sum \frac{p_{iT} \sin\phi_i}{Q} \right| + \mathcal{O}(\frac{p_{iT}^2}{Q^2})$, where $\Delta\phi = \phi_1 - \phi_2$ is the azimuthal angle between the two leptons which have transverse momenta $\vec{p}_{iT}$, $i = 1, 2$, and θ^* is the scattering angle of the dilepton system relative to the beam direction when one boosts to the frame along the beam direction such that the leptons are back to back, one sees that these ϕ_η^*-comparisons also show the underlying physical precision error associated with the defining formula in (10). As expected, the comparisons are somewhat better than the Q_T spectra comparisons because this ϕ_η^* quantity is more inclusive. Two different values of Q_T with correspondingly compensating differences in the attendant ϕ_i can have the same value of ϕ_η^*. We want to encourage the LHC experimentalists to continue to produce the real Q_T spectra in the $Z/\gamma*$ production because we believe, even if there are some irreducible experimental systematic errors, these spectra can be very useful in determining which theoretical approach is actually correct as we argue in the next section. In particular, given the intrinsic physical precision error in (10), one has to account for that error when the data are normalized to any prediction which is based on the formula therein. We would like to stress this latter point. We await the more complete data sets accordingly.

From the discussions just completed, we see that, in order to have a strict control on the theoretical precision in (1), we need both the resummation of the reduced cross section and that of the attendant evolution of the $\{F_j\}$. We turn now to the latter.

More specifically, we apply the QCD restriction of the formula in (4) to the calculation of the kernels, P_{AB}, in the DGLAP-CS theory itself and thereby get an improvement of the IR limit of these kernels, an IR-improved DGLAP-CS theory [7,8] in which large IR effects are resummed for the kernels themselves. The attendant new resummed kernels, $P_{AB}^{\exp}$ are given in Refs. [7–9]. We reproduce the new kernels here for completeness:

$$P_{qq}^{\exp}(z) = C_F F_{\mathrm{YFS}}(\gamma_q) e^{\frac{1}{2}\delta_q} \left[\frac{1+z^2}{1-z}(1-z)^{\gamma_q} - f_q(\gamma_q)\delta(1-z) \right],$$

$$P_{Gq}^{\exp}(z) = C_F F_{\mathrm{YFS}}(\gamma_q) e^{\frac{1}{2}\delta_q} \frac{1+(1-z)^2}{z} z^{\gamma_q},$$

$$P_{GG}^{\exp}(z) = 2C_G F_{\mathrm{YFS}}(\gamma_G) e^{\frac{1}{2}\delta_G} \left\{ \frac{1-z}{z} z^{\gamma_G} + \frac{z}{1-z}(1-z)^{\gamma_G} \right.$$
$$\left. + \frac{1}{2}(z^{1+\gamma_G}(1-z) + z(1-z)^{1+\gamma_G}) - f_G(\gamma_G)\delta(1-z) \right\},$$

$$P_{qG}^{\exp}(z) = F_{\mathrm{YFS}}(\gamma_G) e^{\frac{1}{2}\delta_G} \frac{1}{2}\{z^2(1-z)^{\gamma_G} + (1-z)^2 z^{\gamma_G}\}, \tag{21}$$

where the superscript "exp" indicates that the kernel has been resummed as predicted by Eq. (4) when it is restricted to QCD alone. Here C_F (C_G) is the quadratic Casimir invariant for the quark(gluon) color representation respectively, and the YFS [19] infrared factor is given by

$$F_{\mathrm{YFS}}(a) = e^{-C_E a}/\Gamma(1+a)$$

where $\Gamma(w)$ is Euler's gamma function and $C_E = 0.57721566\ldots$ is Euler's constant. The definitions of the respective resummation functions $\gamma_A, \delta_A, f_A, A = q, G$ are as follows [7,8][5]:

$$\gamma_q = C_F \frac{\alpha_s}{\pi} t = \frac{4C_F}{\beta_0}, \qquad \delta_q = \frac{\gamma_q}{2} + \frac{\alpha_s C_F}{\pi}\left(\frac{\pi^2}{3} - \frac{1}{2}\right),$$

$$f_q(\gamma_q) = \frac{2}{\gamma_q} - \frac{2}{\gamma_q+1} + \frac{1}{\gamma_q+2},$$

$$\gamma_G = C_G \frac{\alpha_s}{\pi} t = \frac{4C_G}{\beta_0}, \qquad \delta_G = \frac{\gamma_G}{2} + \frac{\alpha_s C_G}{\pi}\left(\frac{\pi^2}{3} - \frac{1}{2}\right),$$

$$f_G(\gamma_G) = \frac{n_f}{6C_G F_{\mathrm{YFS}}(\gamma_G)} e^{-\frac{1}{2}\delta_G} + \frac{2}{\gamma_G(1+\gamma_G)(2+\gamma_G)} + \frac{1}{(1+\gamma_G)(2+\gamma_G)}$$
$$+ \frac{1}{2(3+\gamma_G)(4+\gamma_G)} + \frac{1}{(2+\gamma_G)(3+\gamma_G)(4+\gamma_G)}. \tag{22}$$

We use a one-loop formula for $\alpha_s(Q)$ in (22), so that

$$\beta_0 = 11 - \frac{2}{3}n_f \equiv 4\beta_1,$$

where $n_f \equiv N_F$ is the number of active(effective) quark flavors.[6] These new kernels provide us with a new resummed scheme for the parton density functions (PDF's) and the reduced cross section with the same value of σ in (1):

$$F_j, \ \hat{\sigma} \to F_j', \ \hat{\sigma}' \quad \text{for}$$
$$P_{Gq}(z) \to P_{Gq}^{\exp}(z), \quad \text{etc.} \tag{23}$$

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[5] The improvement in Eq. (21) should be distinguished from the resummation in parton density evolution for the "$z \to 0$" Regge regime—see for example Refs. [39,40]. This latter improvement must also be taken into account for precision LHC predictions.

[6] We follow the development of Field in Ref. [41] in removing the mass singularities such that our evolution variable for the DGLAP-CS equation is $t = \ln(Q^2/\Lambda^2)$ as it is given in Eq. (4.6.11) in Ref. [41].
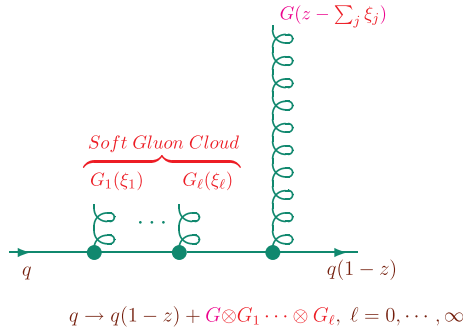
Fig. 1. Bloch–Nordsieck soft quanta for an accelerated charge.

This new scheme has improved MC stability as discussed in Refs. [9]—in the attendant parton shower MC based on the new kernels there is no need for an IR cut-off 'k_0' parameter. It is important to note that, while the degrees of freedom below the IR cut-offs in the usual showers are dropped in those showers, in the showers in HERWIRI1.031, as one can see from (4), these degrees of freedom are included in the calculation and are integrated over in the process of generating the Gribov–Lipatov exponents γ_A in (21). We note also that *the new kernels agree with the usual kernels at $\mathcal{O}(\alpha_s)$ as the differences between them start in $\mathcal{O}(\alpha_s^2)$.* This means that the NLO matching formulas in the MC@NLO and POWHEG frameworks apply directly to the new kernels for the realization of exact NLO ME/shower matching.

In the interest of pedagogy, in Fig. 1 we illustrate the basic physical idea, already discussed by Bloch and Nordsieck [42], which underlies the new kernels: an accelerated charge generates a coherent state of very soft massless quanta of the respective gauge field so that one cannot know which of the infinity of possible states one has made in the splitting process $q(1) \rightarrow q(1-z)+G \otimes G_1 \cdots \otimes G_\ell,\ \ell = 0, \ldots, \infty$ illustrated in Fig. 1. In the new kernels this effect is taken into account by resumming the terms $\mathcal{O}\left((\alpha_s \ln \left(\frac{q^2}{\Lambda^2} \right) \ln(1-z))^n \right)$ when $z \rightarrow 1$ is the IR limit. From (23) and (1), we see that when the usual kernels are used these terms are generated order-by-order in the solution for the cross section σ in (1) so that our resumming them enhances the convergence of the representation in (1) for a given order of exactness in the input perturbative components therein. In the next section, we illustrate this last remark in the context of the comparison of recent LHC data to NLO parton shower/matrix element matched predictions.

3. Interplay of IR-improved DGLAP-CS theory and NLO shower/ME precision: comparison with LHC data

In the new MC HERWIRI1.031 [9] we have the first realization of the new IR-improved kernels in the HERWIG6.5 [10] environment. Here, using recent LHC data as our baseline, we compare it with HERWIG6.510, both with and without the MC@NLO [11] exact $\mathcal{O}(\alpha_s)$ correction to illustrate the interplay between the attendant precision in NLO ME matched parton shower MC's and the new IR-improvement for the kernels.

More specifically, in Fig. 2 in panel (a) we show for the single $Z/\gamma*$ production at the LHC the comparison between the CMS rapidity data [43] and the MC theory predictions and in panel (b) in the same figure we show the analogous comparison with the ATLAS P_T data. Here, the rapidity data are the combined $e^+e^- - \mu^-\mu^+$ results and the p_T data are those for the bare e^+e^- case; for the theoretical framework of our simulations corresponds to these data. We do not as yet have complete realization of all the corrections involved in the other ATLAS data in Ref. [44]. These results are better appreciated if they are considered from the perspective of our analysis in Ref. [9] of the FNAL data on the single Z/γ^* production in p$\bar{\text{p}}$ collisions at 1.96 TeV.

More precisely, we direct the reader to the results in Fig. 11 of the second paper in Ref. [9]. In that figure, we showed that, when the intrinsic rms p_T parameter PTRMS is set to 0 in HERWIG6.5, the
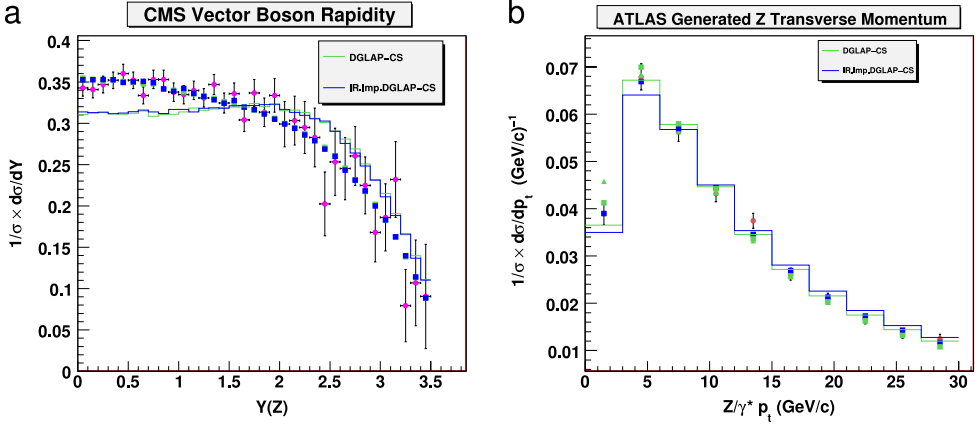
Fig. 2. Comparison with LHC data: (a), CMS rapidity data on (Z/γ^*) production to e^+e^-, $\mu^+\mu^-$ pairs, the circular dots are the data, the green(blue) lines are HERWIG6.510(HERWIRI1.031); (b), ATLAS p_T spectrum data on (Z/γ^*) production to (bare) e^+e^- pairs, the circular dots are the data, the blue(green) lines are HERWIRI1.031(HERWIG6.510). In both (a) and (b) the blue(green) squares are MC@NLO/HERWIRI1.031(HERWIG6.510(PTRMS = 2.2 GeV)). In (b), the green triangles are MC@NLO/HERWIG6.510(PTRMS = 0). These are otherwise untuned theoretical results. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

MC@NLO/HERWIG6.510 simulations give a good fit to the CDF rapidity distribution data [45] therein but they do not give a satisfactory fit to the D0 p_T distribution data [46] therein. In contrast, the corresponding simulations for MC@NLO/HERWIRI1.031 give good fits to both sets of data with the PTRMS = 0. Here PTRMS corresponds to rms value for an intrinsic Gaussian distribution in p_T. The authors of HERWIG [47] already have observed that, to get good fits to both sets of data, one may set PTRMS $\cong$ 2 GeV. Accordingly, in analyzing the new LHC data, we have set PTRMS = 2.2 GeV in our HERWIG6.510 simulations while we continue to set PTRMS = 0 in our HERWIRI simulations.

We turn now with this perspective to the results in Fig. 2, where we see a confirmation of the finding of the HERWIG authors. One needs to set PTRMS $\cong$ 2 GeV [48] in the MC@NLO/HERWIG6.510 simulations to get a good fit to both the CMS rapidity data and the ATLAS p_T data. We again see that the MC@NLO/HERWIRI1.031 simulations with PTRMS = 0 at LHC give a good fit to the data for both the rapidity and the p_T spectra. Quantitatively, we use the χ^2/d.o.f. as a measure of the goodness of the respective fits. From the results in Fig. 2 we compute that the χ^2/d.o.f. for the rapidity data and the χ^2/d.o.f. for the p_T data are (0.72, 0.72)((0.70, 1.37)) for the MC@NLO/HERWIRI1.031 (MC@NLO/HERWIG 6.510(PTRMS = 2.2 GeV)) simulations. The corresponding results are (0.70, 2.23) for the MC@NLO/HERWIG6.510(PTRMS = 0) simulations.

To reproduce the LHC data on the p_T distribution of the Z/γ^* in pp collisions the usual DGLAP-CS kernels require the introduction of a *hard* intrinsic Gaussian distribution in p_T inside the proton whereas the IR-improved kernels give in fact a better fit to the data without the introduction of such a hard intrinsic component to the motion of the proton's constituents. The *hardness* of this intrinsic PTRMS is the issue as this quality of it is entirely ad hoc; it is in disagreement with the results of all successful models of the proton wave-function [49], wherein the scale of the corresponding intrinsic PTRMS is found to be $\lesssim$0.4 GeV. More significantly, it contradicts the well-known experimental observation of *precocious* Bjorken scaling [34,35]; for the famous SLAC-MIT experiments on the deep inelastic electron–proton scattering process show that Bjorken scaling occurs already at $Q^2 = 1_+$ GeV^2 for $Q^2 = -q^2$ with q the 4-momentum transfer from the electron to the proton. If the proton constituents really had a Gaussian intrinsic p_T distribution with PTRMS $\cong$ 2 GeV, these pioneering SLAC-MIT observations would not be possible. What we advocate now is that the ad hoc "hardness" of the PTRMS $\cong$ 2.2 GeV value is really just a phenomenological representation of the more fundamental dynamics described by the IR-improved DGLAP-CS theory. At the level of the inner workings of the two MC's, to understand the physical origin of the difference between the p_T spectra for HERWIRI1.031 and
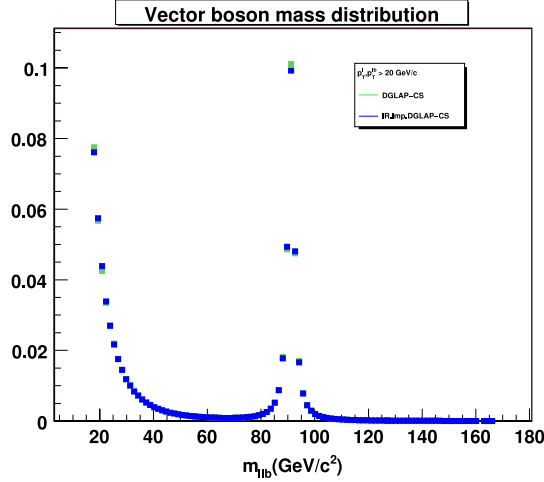
Fig. 3. Normalized vector boson mass spectrum at the LHC for p_T(lepton) > 20 GeV. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

HERWIG6.510 with PTRMS $= 0$, we direct the reader to Fig. 2 in the second paper in Ref. [9]. We see there that the non-branching probability for the soft virtualities is higher for the IR-improved shower in HERWIRI1.031 (blue in the figure) compared to the non-branching probability for the HERWIG6.510 shower with PTRMS $= 0$ (green in the figure) so that there is less relative radiation for p_T near 0 in HERWIRI1.031 shower compared to that in the HERWIG6.510 shower with PTRMS $= 0$, just as we see here in Fig. 2(b).[7] This raises the following question: "Is possible to tell the difference between the two representations of the data in Fig. 2?"

One expects physically that more detailed observations should be able to distinguish the two representations of the data in Fig. 2. In this connection, in Fig. 3 we show for the Z/γ^* mass spectrum the MC@NLO/HERWIRI1.031(blue squares) and MC@NLO/HER-WIG6.510(PTRMS $= 2.2$ GeV) (green squares) predictions when the decay lepton pairs satisfy the LHC type requirement that their transverse momenta $\{p_T^\ell, p_T^{\bar\ell}\}$ exceed 20 GeV. From the results in Fig. 3, wherein the peaks differ by 2.2% for example, we see that the high precision data such as the LHC ATLAS and CMS experiments will have (each already has over 5×10^6 lepton pairs) would allow one to distinguish between the two sets of theoretical predictions.

Continuing in this direction with the discussion, we see that the main differences between the three predictions in Fig. 2(b) occur in the regime below 10 GeV/c. To better probe this latter regime, we make a more detailed snap-shot of it in which we plot in Fig. 4 the three respective featured theory predictions with the finer binning of 0.5 GeV/c instead of the 3.0 GeV/c binning used in Fig. 2(b). The results in Fig. 4 show that the three theoretical predictions have significant differences in the shapes that are testable with the precise data that will be available to the ATLAS and CMS experiments.[8] We would again encourage experimentalists to pursue the measurements of both p_T and ϕ_η^* spectra as these will both be very useful in establishing the correct theoretical approach to the respective LHC

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[7] In Ref. [50] it is shown that the p_T spectra in single Z production at high energy hadron colliders can be explained with a reasonable χ^2/d.o.f. if an ad hoc p_T dependent form of $\alpha_s(p_T)$ at low values of $p_T \to 0$ is assumed; however, when the realistic HERWIG++ hadronization model is used the best-fit values of the two parameter, quadratic ansatz for $\alpha_s(p_T)$ in the non-perturbative regime lead to a parametrization of α_s that is much larger [50] in the non-perturbative regime than those in Refs. [51,52].

[8] We do note that the finer binning we seek for p_T is already available for the new ϕ_η^* variable [53] but that there is significant input (in the context of our precision goals) to the low ϕ_η^* regime from non-soft values of p_T due to the inclusive nature of ϕ_η^* as we have observed.
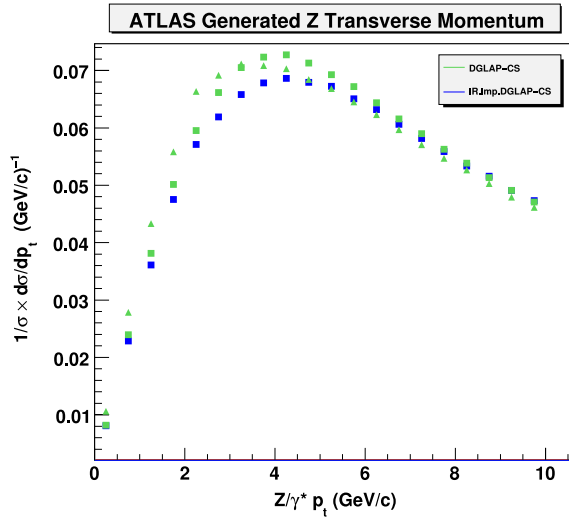
Fig. 4. Normalized vector boson p_T spectrum at the LHC for the ATLAS cuts as exhibited in Fig. 2 for the same conventions on the notation for the theoretical results with the vector boson $p_T < 10$ GeV to illustrate the differences between the three predictions.

observations. We will pursue elsewhere [12] other such detailed observations that may also reveal the differences between the two descriptions of parton shower physics. In this connection especially, we continue to await the release of the entire data sets from ATLAS and CMS.

4. Conclusions

What we have shown is the following. The realization of IR-improved DGLAP-CS theory in HERWIRI1.031, when used in the MC@NLO/HERWIRI1.031 exact $\mathcal{O}(\alpha_s)$ ME matched parton shower framework, affords one the opportunity to explain, on an event-by-event basis, both the rapidity and the p_T spectra of the Z/γ^* in pp collisions in the recent LHC data from CMS and ATLAS, respectively, without the need of an unexpectedly hard intrinsic Gaussian p_T distribution with rms value of PTRMS $\cong 2$ GeV in the proton's wave function. Our view is that this can be interpreted as providing a rigorous basis for the phenomenological correctness of such unexpectedly hard distributions insofar as describing these data using the usual unimproved DGLAP-CS showers is concerned. Accordingly, we have proposed that comparison of other distributions such as the invariant mass distribution with the appropriate cuts and the more detailed Z/γ^* p_T spectra in the regime below 10.0 GeV be used to differentiate between these phenomenological representations of parton shower physics in MC@NLO/HERWIG6.510 and the fundamental description of the parton shower physics in MC@NLO/HERWIRI1.031. We have further emphasized that the precociousness of Bjorken scaling argues against the fundamental correctness of the *hard* scale intrinsic p_T ansatz with the unexpectedly hard value of PTRMS $\cong 2$ GeV, as do the successful models [49] of the proton's wave function, which would predict this value to be $\lesssim 0.4$ GeV. As an added bonus, we have pointed out that the fundamental description in MC@NLO/HERWIRI1.031 can be systematically improved to the NNLO parton shower/ME matched level—a level which we anticipate is a key ingredient in achieving the (sub-)1% precision tag for such processes as single heavy gauge boson production at the LHC.

Evidently, relative to what one could achieve from the fundamental representation of the corresponding physics via IR-improved DGLAP-CS theory as it is realized in HERWIRI1.031 when employed in MC@NLO/HERWIRI1.031 simulations, the use of ad hoc hard scales in models would compromise any discussion of the attendant theoretical precision. We are pursuing additional cross checks of the MC@NLO/HERWIRI1.031 simulations against the LHC data.

In our discussion, we have also spent some amount of time discussing alternative approaches to the type of resummation embodied in the IR-improved DGLAP-CS theory; for some of these approaches are in wide use. What we conclude is that the physical precisions of these approaches (see Section 2), which are based on Refs. [23–26], are above the 1% precision tag to which we aspire, even though there is no contradiction between our exact approach and these more approximate methods.

Acknowledgments

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