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Electroweak corrections to three-jet production in e^+e^- annihilation

Ansgar Denner [a], Stefan Dittmaier [b,c], Thomas Gehrmann [d,*], Christian Kurz [a,d]

[a] Paul Scherrer Institut, CH-5232 Villigen PSI, Switzerland
[b] Physikalisches Institut, Albert-Ludwigs-Universität Freiburg, D-79104 Freiburg, Germany
[c] Max-Planck-Institut für Physik (Werner-Heisenberg-Institut), D-80805 München, Germany
[d] Institut für Theoretische Physik, Universität Zürich, CH-8057 Zürich, Switzerland

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ABSTRACT

We compute the electroweak $\mathcal{O}(\alpha^3\alpha_s)$ corrections to three-jet production and related event-shape observables at electron–positron colliders. We properly account for the experimental photon isolation criteria and for the corrections to the total hadronic cross section. Corrections to the three-jet rate and to normalised event-shape distributions turn out to be at the few-percent level.

Precision QCD studies at electron–positron colliders rely on the measurement of the three-jet production cross section and related event-shape observables. The deviation from simple two-jet configurations is proportional to the strong coupling constant α_s, so that by comparing the measured three-jet rate and related event shapes (see, e.g., Ref. [1]) with the theoretical predictions, one can determine α_s. Including electroweak coupling factors, the leading-order (LO) contribution to this process is of order $\alpha^2\alpha_s$.

Owing to recent calculational progress, the QCD predictions for event shapes [2,3] and three-jet production [4,5] are accurate to next-to-next-to-leading order (NNLO, $\alpha^2\alpha_s^3$) in QCD perturbation theory. Depending on the observable under consideration, the numerical magnitude of the NNLO corrections varies between three and twenty percent. Inclusion of these corrections results in an estimated residual uncertainty of the QCD prediction from missing higher orders at the level of below five percent for the event-shape distributions, and below one percent for the three-jet cross section. At this level of theoretical precision, higher-order electroweak effects could be of comparable magnitude. At present, only partial calculations of electroweak corrections to three-jet production and event shapes are available [6], which cannot be compared with experimental data directly. In this work, we present the first calculation of the NLO electroweak $(\alpha^3\alpha_s)$ corrections to three-jet observables in e^+e^- collisions including the quark–antiquark–photon $(q\bar{q}\gamma)$ final states. Note that the QCD corrections to these final states are of the same perturbative order as the genuine electroweak corrections to quark–antiquark–gluon $(q\bar{q}g)$ final states. Since photons produced in association with hadrons can never be fully isolated, both types of corrections have to be taken into account.

Event-shape measurements at LEP usually rely on a standard set of six variables y, defined, for example, in Ref. [7]: thrust T, C-parameter, heavy jet mass ρ, wide and total jet broadenings B_W and B_T, and two-to-three-jet transition parameter in the Durham algorithm Y_3. The experimentally measured event-shape distribution

$$\frac{1}{\sigma_{had}}\frac{d\sigma}{dy}$$

is normalised to the total hadronic cross section. In the perturbative expansion, it turns out to be most appropriate to consider the

* Corresponding author.
E-mail address: thomas.gehrmann@physik.uzh.ch (T. Gehrmann).

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expansion of this ratio, which reads to NNLO in QCD and NLO in the electroweak theory

$$\frac{1}{\sigma_{\rm had}}\frac{{\rm d}\sigma}{{\rm d}y} = \left(\frac{\alpha_{\rm s}}{2\pi}\right)\frac{{\rm d}\bar{A}}{{\rm d}y} + \left(\frac{\alpha_{\rm s}}{2\pi}\right)^2\frac{{\rm d}\bar{B}}{{\rm d}y} + \left(\frac{\alpha_{\rm s}}{2\pi}\right)^3\frac{{\rm d}\bar{C}}{{\rm d}y} + \left(\frac{\alpha}{2\pi}\right)\frac{{\rm d}\delta_\gamma}{{\rm d}y} + \left(\frac{\alpha_{\rm s}}{2\pi}\right)\left(\frac{\alpha}{2\pi}\right)\frac{{\rm d}\delta_{\rm EW}}{{\rm d}y}, \quad (1)$$

where the fact is used that the perturbative expansion of $\sigma_{\rm had}$ starts at order α^2. The calculation of the QCD coefficients $\bar{A}$, $\bar{B}$, and $\bar{C}$ is described in Refs. [2,3]. The LO purely electromagnetic contribution δ_γ arises from tree-level $q\bar{q}\gamma$ final states without a gluon. The NLO electroweak coefficient $\delta_{\rm EW}$ receives contributions from the $\mathcal{O}(\alpha)$ correction to the hadronic cross section,

$$\sigma_{\rm had} = \sigma_0\left[1 + \left(\frac{\alpha}{2\pi}\right)\delta_{\sigma,1}\right], \quad (2)$$

and from the genuine $\mathcal{O}(\alpha^3\alpha_s)$ contribution to the event-shape distribution

$$\frac{1}{\sigma_0}\frac{{\rm d}\sigma}{{\rm d}y} = \left(\frac{\alpha_{\rm s}}{2\pi}\right)\frac{{\rm d}\bar{A}}{{\rm d}y} + \left(\frac{\alpha_{\rm s}}{2\pi}\right)\left(\frac{\alpha}{2\pi}\right)\frac{{\rm d}\delta_{\bar{A}}}{{\rm d}y},$$

such that

$$\frac{{\rm d}\delta_{\rm EW}}{{\rm d}y} = \frac{{\rm d}\delta_{\bar{A}}}{{\rm d}y} - \delta_{\sigma,1}\frac{{\rm d}\bar{A}}{{\rm d}y} \quad (3)$$

yields the full NLO electroweak correction. Both terms are to be evaluated with the same event-selection cuts. As shown in the following, many of the numerically dominant contributions, especially from initial-state radiation, cancel in this difference.

In the experimental measurement of three-jet observables at electron–positron centre-of-mass energy $\sqrt{s}$, several cuts are applied to reduce the contributions from photonic radiation. In our calculation, we apply the criteria used in the ALEPH analysis [7]. Very similar criteria were also applied by the other experiments [8]. Particles contribute to the final state only if they are within the detector acceptance, defined by the production angle relative to the beam direction, $|\cos\theta| < 0.965$. Events are accepted if the reconstructed invariant mass squared s' of the final-state particles is larger than $s_{\rm cut} = 0.81$ s. To reduce the contribution from hard photon radiation, the final-state particles are clustered into jets using the Durham algorithm with resolution parameter $y_{\rm cut} = 0.002$. If one of the resulting jets contains a photon carrying a fraction $z_\gamma > z_{\gamma,{\rm cut}} = 0.9$ of the jet energy, it is considered to be an isolated photon, and the event is discarded. The event-shape variables are then computed in the centre-of-mass frame of the final-state momenta, which can be boosted relative to the e^+e^- centre-of-mass frame, if particles are outside the detector acceptance.

In the computation of the $\mathcal{O}(\alpha)$ corrections to the total hadronic cross section, we include the virtual electroweak corrections to $q\bar{q}$ final states, and the real radiation corrections from $q\bar{q}\gamma$ final states, provided the above event-selection criteria are fulfilled. The corrections to the event-shape distributions receive contributions from the virtual electroweak corrections to the $q\bar{q}g$ final state, the virtual QCD corrections to the $q\bar{q}\gamma$ final state, and from the real radiation $q\bar{q}g\gamma$ final state. To separate the divergent real radiation contributions, we used both the dipole subtraction method [9,10] and phase-space slicing [11], resulting in two independent implementations. Soft singularities are present in the virtual and real corrections. They are regularized dimensionally or with infinitesimal photon and gluon masses, and cancel in the sum. Collinear singularities from photon radiation off the incoming leptons (initial-state radiation, ISR) are only partially cancelled. The left-over collinear ISR singularity is regularized by the electron mass and absorbed into the initial-state radiator function, which we consider either at fixed order, or in a leading-logarithmic (LL) resummation [12]. Owing to the specific nature of the event selection, also photon radiation off the outgoing quarks (final-state radiation, FSR) is only partially cancelled. The left-over FSR singularity arises from the isolated photon definition, which vetoes on photon jets with $z_\gamma > z_{\gamma,{\rm cut}}$. This singularity is absorbed into the photon fragmentation function, which we apply in the fixed-order approach of Ref. [13]. For the non-perturbative contribution to this function, we use the $\mathcal{O}(\alpha)$ two-parameter fit of ALEPH [14]. The fragmentation contribution derived in Ref. [13] is based on phase space slicing and dimensional regularization. We recomputed this contribution using subtraction and mass regularization [10].

The Feynman diagrams for the virtual corrections are generated with FeynArts [15,16]. Using two independent inhouse Mathematica routines, one of which builds upon FormCalc [17], each diagram is expressed in terms of standard matrix elements and coefficients of tensor integrals. The tensor integral coefficients are numerically reduced to standard scalar integrals using the methods described in Refs. [19,20]. The scalar master integrals are evaluated using the methods and results of Refs. [21–23], where UV divergences are regularized dimensionally. For IR divergences two alternative regularizations are employed, one that is fully based on dimensional regularization with massless light fermions, gluons, and photons, and another that is based on infinitesimal photon and gluon masses and small fermion masses. The loop integrals are translated from one scheme to the other as described in Ref. [24].

The Z-boson resonance is described in the complex-mass scheme [25,26], and its mass is fixed from the complex pole. The electromagnetic couplings appearing in LO are parametrized in the G_μ scheme, i.e., they are fixed via

$$\alpha = \alpha_{G_\mu} = \sqrt{2}G_\mu M_{\rm W}^2\big(1 - M_{\rm W}^2/M_{\rm Z}^2\big)/\pi\,.$$

As the leading electromagnetic corrections are related to the emission of real photons, we fix the electromagnetic coupling appearing in the relative corrections by $\alpha = \alpha(0)$, which is the appropriate choice for the leading photonic corrections. Accordingly the cross section for $e^+e^- \to q\bar{q}g$ is proportional to $\alpha_{G_\mu}^2\alpha_{\rm s}$ while the electroweak corrections to this process are proportional to $\alpha(0)\alpha_{G_\mu}^2\alpha_{\rm s}$.

We performed two independent calculations of all ingredients resulting in two independent Fortran codes, one of them being an extension of Pole [18].

We use the following values of the electroweak and QCD parameters:

$$\begin{aligned} &G_\mu = 1.16637\times 10^{-5}\ {\rm GeV}^{-2}, && M_{\rm H} = 120\ {\rm GeV},\\ &\alpha(0) = 1/137.03599911, && \alpha_{\rm s}(M_{\rm Z}) = 0.1176,\\ &m_{\rm e} = 0.51099892\ {\rm MeV}, && m_{\rm t} = 171.0\ {\rm GeV}. \end{aligned} \quad (4)$$

Because we employ a fixed width in the resonant W- and Z-boson propagators in contrast to the approach used at LEP to fit the W and Z resonances, where running widths are taken, we have to convert the “on-shell” values of $M_V^{\rm LEP}$ and $\Gamma_V^{\rm LEP}$ ($V = {\rm W, Z}$), resulting from LEP, to the “pole values” denoted by M_V and Γ_V, leading to [27]:

$$\begin{aligned} &M_{\rm W} = 80.375\ldots\ {\rm GeV}, && \Gamma_{\rm W} = 2.140\ldots\ {\rm GeV},\\ &M_{\rm Z} = 91.1535\ldots\ {\rm GeV}, && \Gamma_{\rm Z} = 2.4943\ldots\ {\rm GeV}. \end{aligned} \quad (5)$$

In the final state we take all light quarks into account, including b quarks.