NLO electroweak contributions to gluino pair production at hadron colliders

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Abstract
We calculate the full $\mathcal{O}(\alpha_s^2\alpha)$ corrections to the process of gluino pair production at hadron colliders in the framework of the real MSSM. We show that these contributions can be neglected at the LHC performing a scan over a wide region of the parameter space. The impact of these corrections in the parameter range investigated at the Tevatron is small.

1 Introduction

Despite of its success, the Standard Model (SM) of particle physics is affected by theoretical and phenomenological problems whose solution can be used as a guideline for its extension. In this respect Supersymmetry (SUSY) [1, 2], and in particular the Minimal Supersymmetric extension of the Standard Model (MSSM) [3–5], is one of the most promising scenarios for physics beyond the Standard Model. In SUSY models the breaking of electroweak symmetry is obtained radiatively at a scale which is stabilized by Supersymmetry itself. Moreover, the comparison of MSSM predictions and electroweak precision observables provides an overall fit of data [6, 7] which is at least as good as that obtained using the SM, and even better in the case of specific observables such as $g-2$ of the muon [8, 9].

These features render SUSY an appealing framework, and they can explain the big effort in the hunting for SUSY not only in the past but also in future years. In particular, this search is one of the major goals of the LHC. Experimental studies have shown the possibility of early discovery of SUSY with 1 fb$^{-1}$ of integrated luminosity in the inclusive multijet plus missing $E_T$ channel [10, 11], provided that the masses of squarks and gluinos are not too heavy (e.g. $< 2$ TeV). Interestingly, the 95 % confidence level area of the $(m_0, m_{1/2})$ plane of the Constrained MSSM (CMSSM) lies largely in the region that will be investigated with 1 fb$^{-1}$ at 14 TeV [7].
At hadron colliders, colored particles can be searched for most efficiently. This implies that Supersymmetry could be discovered looking at the production of squarks and gluinos. Among the others, the process of gluino pair production,

\[ P P \rightarrow \tilde{g} \tilde{g} X, \]  

is one of the most important processes leading to the production of colored SUSY particles. Indeed, its cross section is large, \( \mathcal{O}(10 \text{ pb}) \) if the gluino mass is \( \mathcal{O}(600 \text{ GeV}) \). Moreover, the gluino plays a key role in characterizing SUSY models. The measure of the spin of a (supposed to be) gluino \[12\] and the confirmation of its Majorana nature \[13–15\] would allow not only to distinguish among different beyond Standard Model scenarios, but also among MSSM and others SUSY models involving Dirac gauginos, such as the \( N = 1/N = 2 \) hybrid scheme \[16\].

The total cross section for gluino pair production was computed at Born level long time ago \[17–20\]. NLO SUSY-QCD contributions were computed in Ref. \[21\]. These corrections are positive and large (from 5 to 90%, depending on the masses of the squarks and of the gluino), and they reduce appreciably the factorization scale dependence. They are included into the publicly available code \texttt{Prospino} \[22\]. More recent is the resummation of the QCD Sudakov logarithms at the next-to-leading-logarithmic (NLL) accuracy \[23, 24\], and the resummation of the leading Coulomb corrections \[24\]. Their inclusion further stabilizes the prediction against scale variation. The NLL contributions are of the order of 2–8\% of the NLO QCD predictions, provided the squark and gluino masses are \( \mathcal{O}(1 \text{ TeV}) \). In this mass range the contribution of the Coulomb corrections amounts up to 5\%. The computation of NLO electroweak (EW) corrections, of \( \mathcal{O}(\alpha_s^2\alpha) \), to the process of hadronic production of a gluino pair is still missing. In this paper, we fill this gap computing the full \( \mathcal{O}(\alpha_s^2\alpha) \) corrections to the process (1) in the framework of the real MSSM. Our computation is part of an ongoing project aiming to evaluate the tree-level EW and NLO EW contributions to the production of colored SUSY particles at the LHC \[25–29\].

The plan of the paper is the following. In section 2 we briefly summarize the \( \mathcal{O}(\alpha_s^3) \) contributions to the process (1). In section 3 we describe the partonic processes contributing at \( \mathcal{O}(\alpha_s^2\alpha) \). Numerical results for the electroweak corrections to gluino pair production at the LHC are presented in section 4, while section 5 is devoted to a brief discussion on the numerical value of the electroweak corrections at the Tevatron. Section 6 summarizes our results. A list of Feynman diagrams is collected in the appendix.

## 2 Gluino pair production in lowest order

The leading order contributions to the process (1) are of QCD origin, of \( \mathcal{O}(\alpha_s^3) \). At lowest order in the perturbative expansion the differential cross section can be written
as follows,
\[ d\sigma_{pp→\bar{g}g}^{\text{cd}, \text{LO}}(S) = \sum_q \int_{\tau_0}^1 d\tau \frac{dL_{pq}}{d\tau}(\tau)d\sigma_{pq→g\bar{g}}^{2,0}(\tau S) + \int_{\tau_0}^1 d\tau \frac{dL_{gg}}{d\tau}(\tau)d\sigma_{gg→g\bar{g}}^{2,0}(\tau S), \]

(2)

with the help of the parton luminosities, defined according to
\[ \frac{dL_{ij}}{d\tau}(\tau) = \frac{1}{1+\delta_{ij}} \int_\tau^1 \frac{dx}{x} \left[ f_{ij}(x)f_{j|i}(\frac{\tau}{x}) + f_{ji}(x)f_{ij}(\frac{\tau}{x}) \right]. \]

(3)

where \( f_{iH}(x) \) is the momentum distribution of the parton \( i \) inside the hadron \( H \). The sum runs over the quarks \( q = u, d, c, s, b \). The lower limit on the integral over \( \tau \), \( \tau_0 = 4 m_\tilde{g}^2 / S \), is related to the threshold for the production of the gluino pair. We use the convention \( d\sigma_X^{ab} \) to denote the cross section for a partonic process \( X \) at a given order \( O(\alpha_s^a\alpha_s^b) \) in the strong and electroweak coupling constants. Therefore, \( d\sigma_{qq→g\bar{g}}^{2,0} \) and \( d\sigma_{gg→g\bar{g}}^{2,0} \) are the lowest order differential cross sections for the partonic processes

\[ q(p_1)\bar{q}(p_2) \rightarrow \tilde{g}(k_1)\tilde{g}(k_2), \]

(4)

\[ g(p_1)g(p_2) \rightarrow \tilde{g}(k_1)\tilde{g}(k_2), \]

(5)

respectively. The cross sections are averaged (summed) over the spins and the colors of the incoming (outgoing) particles. In this analysis we will consider the five light quarks as massless and we approximate the CKM matrix by the unity matrix. We retain the mass of the bottom in the Yukawa couplings, owing to the possible enhancement due to \( \tan \beta \). We perform our computation in Feynman gauge.

In lowest order, the partonic cross sections for the processes (4) and (5) can be obtained from the Feynman diagrams in Fig. 1 of appendix A. The cross sections can be written as

\[ d\sigma_{qq→g\bar{g}}^{2,0} = \frac{dt}{16\pi s^2} \sum |\mathcal{M}_{qq→g\bar{g}}^{1,0}|^2, \]

\[ d\sigma_{gg→g\bar{g}}^{2,0} = \frac{dt}{16\pi s^2} \sum |\mathcal{M}_{gg→g\bar{g}}^{1,0}|^2, \]

(6)

where \( \mathcal{M}_{X}^{1,0} \) is the tree-level contribution to the amplitude of the process \( X \). The squared amplitudes averaged (summed) over the spins and the colors of the initial (final) particles read [17–21]

\[ \sum |\mathcal{M}_{qq→g\bar{g}}^{1,0}|^2 = 8 \alpha_s^2 \pi^2 \left\{ \frac{72}{s^2} \left( 2m_\tilde{g}^2 s + t_\tilde{g}^2 u_\tilde{g} \right) + 4m_\tilde{g}^2 s \left( \frac{1}{u_{\tilde{q},1}^2 t_{\tilde{q},1}} + \frac{1}{u_{\tilde{q},2} t_{\tilde{q},2}} \right) + 36(m_\tilde{g}^2 s + t_\tilde{g}^2) \left( \frac{1}{t_{\tilde{q},1}} + \frac{1}{t_{\tilde{q},2}} \right) + 16t_\tilde{g}^2 \left( \frac{1}{t_{\tilde{q},1}^2} + \frac{1}{t_{\tilde{q},2}^2} \right) \right\}, \]

\[ \sum |\mathcal{M}_{gg→g\bar{g}}^{1,0}|^2 = 18 \alpha_s^2 \pi^2 \left\{ \left( 1 - \frac{t_{\tilde{g}} u_{\tilde{g}}}{s^2} \right) \left[ \frac{s^2}{t_{\tilde{g}} u_{\tilde{g}}} - 2 + 4 \frac{m_\tilde{g}^2 s}{t_{\tilde{g}} u_{\tilde{g}}} \left( 1 - \frac{m_\tilde{g}^2 s}{t_{\tilde{g}} u_{\tilde{g}}} \right) \right] \right\}. \]

(7)
A factor $1/2$ has been taken into account because of the identical particles in the final states. The Mandelstam variables are defined as

\begin{align*}
    s &= (p_1 + p_2)^2, \\
    t &= (p_2 - k_2)^2, \quad t_{\tilde{q},a} = t - m_{\tilde{q},a}^2, \quad t_{\tilde{g}} = t - m_{\tilde{g}}^2, \\
    u &= (p_1 - k_2)^2, \quad u_{\tilde{q},a} = u - m_{\tilde{q},a}^2, \quad u_{\tilde{g}} = u - m_{\tilde{g}}^2,
\end{align*}

and $u = 2m_{\tilde{g}}^2 - s - t$.

### 3 $\mathcal{O}(\alpha_s^2 \alpha)$ corrections to the hadronic process

Gluinos do not interact weakly, thus a pair of gluinos is neither produced at $\mathcal{O}(\alpha^2)$, from $q\bar{q}$ initial states, nor at $\mathcal{O}(\alpha_s \alpha)$ via photon-induced processes. Therefore, in contrast to squark–anti-squark [25, 27, 28] and squark-gluino [29] production, EW contributions enter only at NLO, and they are at least of $\mathcal{O}(\alpha_s^2 \alpha)$. The NLO EW contributions to the hadronic differential cross section reads as follows,

\begin{equation}
    d\sigma_{\text{PP} \to \tilde{g}\tilde{g}X}^{\text{NLO}}(S) = \sum_q \left\{ \int_{\tau_0}^{1} d\tau \left[ \frac{dL_{q\bar{q}}}{d\tau}(\tau) \left( d\sigma_{q\bar{q} \to \tilde{g}\tilde{g}}^{2,1}(\tau S) + d\sigma_{q\bar{q} \to \tilde{g}\tilde{g}\gamma}^{2,1}(\tau S) \right) + \frac{dL_{q\gamma}}{d\tau}(\tau) d\sigma_{q\gamma \to \tilde{g}\tilde{g}}^{2,1}(\tau S) \right] \right\}. \tag{9}
\end{equation}

The $q\gamma$ and $\gamma q$ luminosities entering (9) are built according to Eq. (3). Besides the virtual corrections and the real photon radiation processes at $\mathcal{O}(\alpha_s^2 \alpha)$, we consider the photon-induced processes leading to the production of a gluino pair together with an (anti-)quark. Diagrams and amplitudes are generated with FeynArts [30, 31]. The reduction of the one-loop integrals is performed with the help of FormCalc [32, 33], while the scalar one-loop integrals are numerically evaluated using LoopTools. Infrared (IR) and collinear singularities are treated using mass regularization, i.e. giving a small mass to the photon and to the five light quarks.

#### 3.1 $q\bar{q}$ annihilation with electroweak loops

The first class of corrections entering Eq. (9) are the electroweak one-loop corrections to the processes (4), yielding the following partonic cross section,

\begin{equation}
    d\sigma_{q\bar{q} \to \tilde{g}\tilde{g}}^{2,1} = \frac{dt}{16\pi s^2} \sum_q 2 \text{Re} \left\{ \mathcal{M}_{q\bar{q} \to \tilde{g}\tilde{g}}^{1,0} \mathcal{M}_{q\bar{q} \to \tilde{g}\tilde{g}}^{1,1} \right\}. \tag{10}
\end{equation}

$\mathcal{M}_{q\bar{q} \to \tilde{g}\tilde{g}}^{1,1}$ is the one-loop electroweak contribution to the amplitude of the quark–anti-quark annihilation process. The diagrams responsible for this contribution are displayed in Fig. 4 of appendix A.
We treat UV divergences using dimensional reduction. In order to cure the UV divergences we have to renormalize the quark and the squark sector at $O(\alpha)$. Renormalization of mass and wavefunction of the quarks and squarks belonging to the first two generations, of $\tan \beta$ and of the mass of the $W$ boson, has been performed according to the procedure described in Ref. [28]. The renormalization of mass and wavefunction of the bottom and top quarks and squarks has been widely studied [34–38], and several renormalization schemes have been proposed. Each of these schemes has its own virtues and drawbacks, we perform our computation using two different renormalization schemes. The first (second) scheme, referred in the following as Rs1 (Rs2), is the ”$m_b$ OS” (”$m_b$ DR”) scheme defined in Ref. [38]. The renormalization of the stop-bottom sector at $O(\alpha)$ within the Rs2 scheme requires the renormalization of the supersymmetric Higgs parameter $\mu$. This parameter has been defined in the DR scheme.

In the case of $b\bar{b} \rightarrow \tilde{g}\tilde{g}$ we keep the mass of the $b$-quark that appears in the couplings. In this case, the last twelve diagrams in Fig. 4 of the appendix A have to be considered as well. It is well known [39–44] that, in the large $\tan \beta$ regime, the tree-level relation between the bottom mass $m_b$ and the bottom Yukawa couplings $y_b$ receives radiative corrections that can be strongly enhanced and have to be resummed. Power counting in $\alpha_s\tan \beta$ shows that the leading $\tan \beta$ enhanced contributions, of $O(\alpha_s^2 \tan^n \beta)$, can be accounted for by means of the substitution

$$m_b^{Rs} \rightarrow \overline{m}_b^{Rs} = \frac{m_b^{Rs}}{1 + \Delta_b} \quad (11)$$

in the relation between $m_b$ and $y_b$. $m_b^{Rs}$ is the bottom mass in a given renormalization scheme, Rs. $\Delta_b$ is defined as

$$\Delta_b = \frac{2\alpha_s}{3\pi} \frac{m_\tilde{g} \mu \tan \beta}{(m_{b,1}^2 - m_{b,2}^2)(m_{b,2}^2 - m_\tilde{g}^2)(m_{b,1}^2 - m_{b,2}^2)} \left[ m_{b,1}^2 m_{b,2}^2 \ln \left( \frac{m_{b,1}^2}{m_{b,2}^2} \right) \right. + \left. m_{b,2}^2 m_\tilde{g}^2 \ln \left( \frac{m_{b,2}^2}{m_\tilde{g}^2} \right) + m_\tilde{g}^2 m_{b,1}^2 \ln \left( \frac{m_\tilde{g}^2}{m_{b,1}^2} \right) \right]. \quad (12)$$

Concerning the Higgs sector, the $b\bar{b} - H_0$ coupling is dynamically generated at $O(\alpha_s)$. This coupling can be enhanced if $\tan \beta$ is large and it is worth to include such effects modifying the $b\bar{b}$–Higgs Yukawa couplings. In particular, the effective Lagrangian that correctly takes into account these dynamically generated extra-couplings is

$$L_{\text{eff, Higgs}}^{\text{string}} = \frac{\pi^{Rs}}{v} \left[ \tan \beta \left( 1 - \frac{\Delta_b}{\tan^2 \beta} \right) A^0 \tilde{b}i\gamma^5 \tilde{b} + \frac{\sin \alpha}{\cos \beta} \left( 1 - \frac{\Delta_b}{\tan \alpha \tan \beta} \right) H^0 \bar{b}b \right. + \left. \frac{\cos \alpha}{\cos \beta} \left( 1 + \frac{\Delta_b}{\tan \alpha \tan \beta} \right) H^0 \bar{b}b \right]. \quad (13)$$
3.2 \( q\bar{q} \) annihilation with real photon emission

IR singularities in the virtual corrections are cancelled when the tree-level contribution of the partonic process of real photon radiation,

\[
q(p_1) \bar{q}(p_2) \rightarrow \tilde{g}(k_1) \tilde{g}(k_2) \gamma(k_3),
\]

is included. This contribution can be computed using the Feynman diagrams depicted in Fig. 2 of appendix A. The integral over the phase space is IR divergent when \( k_3 \rightarrow 0 \), while collinear singularities appear whenever \( k_3 \cdot p_i \rightarrow 0 \). IR and collinear divergences appearing in the phase space integration are regularized using both, phase space slicing [45–47] and dipole subtraction [48, 49]. The two methods are in good numerical agreement, as found also in the case of squark-anti–squark production [28]. As already mentioned, IR singularities cancel when the real radiation processes and the virtual contributions are added together, as in Eq. (9). Collinear singularities remain and have to be absorbed via the factorization of the parton distribution functions (PDFs), c.f. section 3.4.

3.3 \( q\gamma \) and \( \gamma\bar{q} \) fusion

The last class of \( \mathcal{O}(\alpha_s^2 \alpha) \) contributions to the process (1) are the tree-level contributions of the partonic processes

\[
\begin{align*}
q(p_1) \gamma(p_2) & \rightarrow \tilde{g}(k_1) \tilde{g}(k_2) q(k_3), \\
\gamma(p_1) \bar{q}(p_2) & \rightarrow \tilde{g}(k_1) \tilde{g}(k_2) \bar{q}(k_3).
\end{align*}
\]

These contributions can be computed from the Feynman diagrams depicted in Fig. 3 of appendix A.

Note that, if \( m_{\tilde{q}} > m_{\tilde{g}} \), the quark in the final state can be the decay product of an on-shell squark. If this is the case the last four diagrams depicted in Fig. 3 become singular. The related poles have to be regularized inserting the width of the on-shell squarks into the corresponding propagator. Furthermore, the contribution obtained squaring the resonant diagrams has to be subtracted since it arises from the production and the subsequent decay of an (anti-)squark through (anti-)quark–photon fusion,

\[
\begin{align*}
q \gamma & \rightarrow \tilde{g}\tilde{q} \quad \text{and} \quad \bar{q} \rightarrow \tilde{g} q, \\
\gamma \bar{q} & \rightarrow \tilde{g}\tilde{q}^* \quad \text{and} \quad \bar{q}^* \rightarrow \tilde{g} \bar{q}.
\end{align*}
\]

According to Refs. [21, 29], the extraction of the Breit-Wigner pole contribution has been performed in the narrow width approximation.

Collinear singularities arising from initial state emission are again absorbed into the PDFs. These singularities are regularized using both, phase space slicing and dipole subtraction. The formulae needed can be found in Ref. [45] and in Refs. [49, 50],
respectively. The results obtained using the two methods agree within the integration uncertainty.

The contribution of this channel is expected to be small owing to the suppression of the photon PDF inside the proton. Indeed, the photon PDF is intrinsically suppressed with respect to the valence quark PDF by a factor $\alpha$, since this PDF is originated from the emission of a photon from a (anti-)quark. In the SUSY scenarios we consider, the contribution of this partonic process amounts up to few percent of the whole $O(\alpha_s^2 \alpha)$ correction.

### 3.4 Factorization of initial collinear singularities

As already mentioned, the universal logarithmic divergences related to the collinear splittings

$$ q \to q \gamma, \quad \bar{q} \to \bar{q} \gamma, \quad \gamma \to q \bar{q}, $$

are absorbed into the definition of the PDFs via mass factorization. We factorize the (anti-)quark PDFs at $O(\alpha)$ in the DIS scheme. The effect of this factorization is to add the following term into Eq. (9),

$$ d\sigma_{\text{Fact}} = \int_{z_0}^1 \frac{d\tau}{\tau} \sum_q \left\{ -\frac{\alpha}{\pi} e_q^2 \frac{dL_{q\gamma}}{d\tau}(\tau) \int_{z_0}^1 dz \left[ H_q^{(1)} \right] \right. + \left. \int_{z_0}^1 dz \left( H_q^{(2)} \right) d\sigma_{q\gamma \to \bar{q}\bar{g}}(z \tau S) \right\} d\sigma_{q\gamma \to \bar{q}\bar{g}}(z \tau S). \tag{18} $$

The functions $H_q^{(1)}$ and $H_q^{(2)}$ read as follows,

$$ H_q^{(1)} = P_{qq}(z) \left[ \ln \left( \frac{\mu^2}{m_q^2(1-z)} - 1 \right) - 3 \frac{1}{2(1-z)} + 2z + 3, \right. \right.$$

$$ \left. \left. H_q^{(2)} = P_{q\gamma}(z) \ln \left( \frac{\mu^2}{m_q^2} \frac{1-z}{z} \right) - 1 + 8z - 8z^2, \right. \right.$$

where the splitting functions are

$$ P_{qq}(z) = \frac{1 + z^2}{1-z}, \quad P_{q\gamma}(z) = z^2 + (1-z)^2. $$

$z_0$ is defined as $z_0 = 4m_{\bar{g}}/(\tau S)$, while $e_q$ is the charge of the quark $q$ expressed in units of the positron charge. The $[\cdots]_+$ distribution is defined as

$$ \int_a^1 dx \left[ f(x) \right]_+ g(x) = \int_a^1 dx \ f(x)[g(x) - g(1)] - g(1) \int_0^a dx \ f(x). \tag{20} $$

In the actual computation, we use the MRST2004qed parton distribution functions at NLO QED and NLO QCD [51]. This fit takes into account QED-effects into
Table 1: MSSM input parameters for the computation of the spectrum of the two scenarios considered. $m_{1/2}$, $m_0$ and $A_0$ are defined at the GUT scale.

| parameter | SPS1a'    | SPS2     |
|-----------|-----------|----------|
| $m_{1/2}$ | 250 GeV   | 300 GeV  |
| $m_0$     | 70 GeV    | 1450 GeV |
| $A_0$     | $-300$ GeV| 0        |
| $\text{sign}(\mu)$ | " +" | " +" |
| $\tan\beta(M_Z)$ | 10.37 | 10 |

the DGLAP evolution equations and the parametrization of the PDF at the initial scale. MRST2004qed PDFs are defined at NLO QCD within the $\overline{\text{MS}}$ mass factorization scheme. As discussed in Ref. [50], the DIS scheme is used for the factorization of the $\mathcal{O}(\alpha)$ corrections.

In our computation we set the renormalization scale, $\mu_R$, equal to the factorization scale, $\mu_F$, and to the gluino mass, i.e. $\mu_R = \mu_F = m_{\tilde{g}}$.

4 Numerical results, LHC

For our numerical discussion we use the Standard Model parameters quoted in Ref. [52]. The value of the bottom mass in the DR scheme is computed according to Ref. [38]. We choose two different SUSY scenarios. The first scenario is the SPS1a' suggested by the Supersymmetry Parameters Analyses (SPA) [53] project. The second one, called SPS2, belongs to the set of Snowmass Points and Slopes, introduced in Ref. [54]. We obtain the parameters of the two scenarios with the help of the program SPheno [55], starting from the input parameters shown in Table 1. In the SPS1a' (SPS2) scenario, the gluino mass is 608 GeV (784 GeV).

4.1 Dependence on the SUSY scenario

We compute the total hadronic cross section, the results are collected in Table 2. The second column shows the lowest order results. The third column shows the sum of the lowest order and of the $\mathcal{O}(\alpha_s^2\alpha)$ contributions. In the fourth column the contribution of the $\mathcal{O}(\alpha_s^2\alpha)$ corrections relative to the total result is given, i.e. $\delta$ is defined as

$$\delta \equiv \frac{\sigma_{\text{EW}, \text{NLO}}}{\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}}. $$

In the last entry we give an estimate of the statistical error based on an integrated luminosity $L = 100$ fb$^{-1}$ [56]. We do not distinguish the results in the different renormalization schemes since they agree within the integration error. A priori this is not guaranteed. Indeed the Rs1 scheme turns out to be unreliable in the scenarios we are
Table 2: Total hadronic cross section for gluino pair production at the LHC ($\sqrt{S} = 14$ TeV). In the second (third) column we show the $\mathcal{O}(\alpha_s^2)$ ($\mathcal{O}(\alpha_s^2 + \alpha_s^2 \alpha)$) contribution for the points SPS1a' and SPS2. In the fourth column the electroweak corrections relative to the LO + NLO EW result are given. The last column shows the statistical error for an integrated luminosity $L = 100$ fb$^{-1}$.

| point     | $\sigma_{QCD, LO}$    | $\sigma_{QCD, LO} + \sigma_{EW, NLO}$ | $\delta$ | $\sqrt{\sigma_{lo, QCD, LO}}$ |
|-----------|-----------------------|----------------------------------------|----------|-------------------------------|
| SPS1a'    | 6.1865(6) pb          | 6.1822(6) pb                          | -0.07%   | 0.13%                         |
| SPS2      | 1.2127(1) pb          | 1.2089(1) pb                          | -0.31%   | 0.29%                         |

considering. In this scheme, the finite part of the renormalization constant of the trilinear coupling, $\delta A_{b}^{fin}$, is comparable with the value of the trilinear coupling $A_b$ itself, i.e. $\delta A_{b}^{fin}/A_b \sim 1$, and the perturbative expansion is spoiled. However, the difference among the renormalization schemes is as small as few percent of the tree-level $b\bar{b}$ annihilation channel cross section. The latter contributions amount up to several fb, therefore the variation of the results in Table 2 is within the integration error.

As one can see, in the case of the point SPS1a' the electroweak corrections are much smaller than the statistical uncertainty and so they are not relevant. In the case of the point SPS2, the $\mathcal{O}(\alpha_s^2 \alpha)$ corrections are of the same order of the statistical error but they are smaller than the theoretical systematic uncertainties such as the uncertainty on the PDF parametrization ($\lesssim 10\%$) [21] and the factorization scale dependence (from 3 to 5% if $m_{\tilde{g}} \leq 1$ TeV) [24].

The invariant mass distribution of the two gluinos is shown in Fig. 5. The EW corrections are small, their absolute value being at most of the order of 0.4% of the total contribution. Moreover, these corrections do not distort the shape of the distribution.

In Fig. 6 we consider the distribution of the largest transverse momentum of the two gluinos, for brevity we will refer to this observable as "transverse momentum distribution". The $\mathcal{O}(\alpha_s^2 \alpha)$ corrections are rather small, the absolute value of their contribution relative to the total result is at most 1%, reaching this value in the high $p_T$ region, for $p_T \gtrsim 1500$ GeV.

4.2 Dependence on the MSSM parameters

In this subsection we investigate the size of the $\mathcal{O}(\alpha_s^2 \alpha)$ corrections to gluino pair production in a more systematic way, performing a scan over the parameter space of the MSSM. The parameters involved in the scan are the independent parameters in the second renormalization scheme. We suppose that all the sfermionic soft mass parameters are equal and we indicate their value in the Rs2 scheme as $M_{Susy}$. The physical masses of the sfermions can be obtained from $M_{Susy}$ diagonalizing the mass
matrices. Moreover, we consider the surfaces of the parameter space characterized by $A_t = A_t$. With these assumptions there are eight independent parameters involved in the scan, namely,

$$M_{\text{Susy}}, \ m_{\tilde{g}}, \ \mu, \ M_2, \ A_t, \ A_b, \ \tan \beta, \ M_{A^0}. \ $$

The subregions of the parameter space are chosen imposing the exclusion limits arising from SUSY searches at LEP [57] and at the Tevatron [58], and the bound on the mass of the light Higgs boson. The physical mass of the light Higgs boson has been computed using *FeynHiggs 2.5.1* [59–61]. Moreover, each point in the selected regions fulfills the condition $|\Delta \rho| \leq 0.025$, $\Delta \rho$ being the dominant SUSY corrections to the electroweak $\rho$ parameter, corrections arising from top and bottom squarks contributions.

We perform four different scans. In each scan we select two parameters and we study the dependence of the quantity $\Delta$,

$$\Delta \equiv \frac{\sigma_{\text{EW, NLO}}}{\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}} \cdot 100. \ $$

We repeat each of these scans for different values of another pair of parameters, while the remaining four are fixed to their SPS1a' values. Here there is a brief discussion on the results of these scans.

**Scan over $A_t$ and $A_b$**

The results of this scan are displayed in Fig. 7. As expected, $\Delta$ is quite independent on the parameter $A_t$ which enters in the virtual correction of the process $b\bar{b} \rightarrow \tilde{g}\tilde{g}$ and in the definition of the mass of the top squarks. This feature is more evident for large $\tan \beta$ values. $\Delta$ varies only by an amount of the order of few percent for a variation of $A_b$ and $A_t$ over a quite broad range (from $-1500$ to $1500$ GeV ). Note that in the whole subregion considered the absolute value of $\Delta$ is of $\mathcal{O}(0.1)$.

**Scan over $\tan \beta$ and $M_{A^0}$**

As can be inferred from Fig. 8, the dependence of $\Delta$ on $(\tan \beta, M_{A^0})$ strongly varies for different values of $m_{\tilde{g}}$ and $M_{\text{Susy}}$. As a general result the overall dependence is mild for each value of $(m_{\tilde{g}}, M_{\text{Susy}})$. In all cases the value of $|\Delta|$ is at most of the order of 0.7.

**Scan over $\mu$ and $M_2$**

As displayed in Fig. 9, $\Delta$ is almost independent on $\mu$ for each value of the pair $(m_{\tilde{g}}, M_{\text{Susy}})$ while the dependence on $M_2$ is more important and particularly pronounced when $m_{\tilde{g}} = 1250$ GeV and $M_{\text{Susy}} = 730$ GeV. In the case of the last three plots the value of $\Delta$ is of order $-0.3$ to $-0.1$, while in the first plot, characterized by $m_{\tilde{g}} \sim 2 \cdot M_{\text{Susy}}$, the value of $\Delta$ is enhanced for small values of $M_2$ reaching the value of $-0.65$. Notice that the mass of the lightest neutralino and chargino is almost independent on
the value of $\mu$ but varies strongly as $M_2$ varies, growing as the value of this parameter grows. So this enhancement occurs when charginos/neutralinos are much lighter than the gluino.

**Scan over $M_{\text{Susy}}$ and $m_{\tilde{g}}$**

In this scan we investigate the dependence of $\Delta$ on $m_{\tilde{g}}$ and $M_{\text{Susy}}$, which is expected to be the most important because of the dependence of the lowest order cross section on these parameters. We consider the variation of $\Delta$ as a function of $(m_{\tilde{g}}, M_{\text{Susy}})$ for different values of $M_2$ and $\tan \beta$, see Fig. 10. Note that we plot $\xi \equiv -\Delta$ instead of $\Delta$. As a general feature $\xi$ increases as $m_{\tilde{g}}$ increases and as $M_{\text{Susy}}$ decreases. The behaviour of $\xi$ as a function of $M_{\text{Susy}}$ and $m_{\tilde{g}}$ is affected by the value of $M_2$ being enhanced for smaller values of this parameter. In particular $\xi \sim 3$ in the region $m_{\tilde{g}} \geq 1600$ GeV, $M_{\text{Susy}} \leq 500$ GeV.

The enhancement of the EW corrections is related to the increasing importance of the $q\bar{q}$ annihilation channel when the production threshold becomes higher. Indeed, the minimal value of the parton’s momentum fraction rises as the gluino mass rises. Since the relative importance of the (anti-)quark PDF increases as the momentum fraction of the (anti-)quark increases, the EW corrections grow as the mass of the gluino grows. The relative importance of the EW contributions is more pronounced when $M_{\text{Susy}}$ is small owing to the presence of tree-level diagrams with squarks exchanged in the $t$ and $u$ channel, c.f. Fig. 1, which are enhanced when the squark masses decrease.

## 5 Numerical results, Tevatron

The EW contributions to gluino pair production are expected to be more important at the Tevatron than at the LHC, owing to the enhancement of the quark–anti-quark annihilation channels with respect to the gluon fusion channel. Therefore, it is worth to estimate the impact of the EW contributions to gluino pair production at the Tevatron, i.e. to the process

$$P \bar{P} \to \tilde{g} \tilde{g} X.$$  \hspace{1cm} (21)

The previous analysis can be easily extended to (21), provided that the definition of the luminosity, Eq. (3), is replaced by

$$\frac{dL_{ij}}{d\tau}(\tau) = \frac{1}{1 + \delta_{ij}} \int_{\tau}^{1} \frac{dx}{x} \left[ f_{ijP}(x)f_{jiP}(\frac{\tau}{x}) + f_{jiP}(\frac{\tau}{x}) f_{ijP}(x) \right].$$  \hspace{1cm} (22)

For numerical evaluation, we focus on two different points of the MSSM parameter space, referred to as TP1 and TP2 respectively. These points belong to the region of the parameter space of the MSSM used in the data analysis made by CDF and D0 collaborations [62–64]. We obtain the parameters in these scenarios with the help of SPheno, starting from the input parameters at the GUT scale described in Table 3.
Table 3: MSSM input parameters for the computation of the spectrum of the scenarios TP1 and TP2.

| parameter | TP1 | TP2 |
|-----------|-----|-----|
| $m_{1/2}$ | 200 GeV | 500 GeV |
| $m_0$     | 130 GeV | 120 GeV |
| $A_0$     | 0     | 0    |
| sign($\mu$) | "−" | "−" |
| $\tan\beta(M_Z)$ | 3 | 3 |

Table 4: Same as Table 2, but considering gluino pair production at the Tevatron, i.e. the process $p\bar{p} \rightarrow \tilde{g}\tilde{g}X$ at $\sqrt{S} = 1.96$ TeV, and different SUSY scenarios. In this case, the last column shows the statistical error for an integrated luminosity $L = 2 \times 8$ fb$^{-1}$.

These points are compatible with the experimental limits set by the analysis made by the D0 collaboration [64]. In particular, the first one corresponds to a scenario in which the gluino is heavier than the squarks ($m_{\tilde{g}} \sim 500$ GeV, $m_{\tilde{q} \neq \tilde{t}} \sim 460$ GeV), while the second one describes a scenario characterized by a light gluino ($m_{\tilde{g}} \sim 340$ GeV, and $m_{\tilde{q} \neq \tilde{t}} \sim 550$ GeV).

In Table 4 we show the total hadronic cross section in the two points considered. We use the same notation as in Table 2. In the case of the point TP1, the size of the electroweak corrections is so small that they will not be visible at the expected final integrated luminosity $L = 2 \times 8$ fb$^{-1}$. In the case of the point TP2 we obtain a relative statistical error of order 4% which is three times bigger than the size of the electroweak contributions. Moreover it is worth to notice that the systematic uncertainties affecting SUSY searches at the Tevatron are typically greater than 1%. For instance, Ref. [64] claims that the $\mu_F$ dependence of the total cross section gives an error from 15 to 20%.

The invariant mass distribution for the two points is shown in Fig. 11. In both cases the $\mathcal{O}(\alpha_s^2\alpha)$ corrections are small compared to the lowest order results and do not change the shape of the distribution. In particular, in the case of the point TP1 (TP2) EW corrections relative to the total contribution are of the order of $-2$ to 5% ($-3$ to $-1$%).

Similar considerations hold in the case of the transverse momentum distribution, shown in Fig. 12. The shape of the distribution is not affected by the insertion of the electroweak corrections in both points. The electroweak corrections relative to the
total contribution are of order of $-2$ to $4.5\%$ in the case of the TP1 point and of the order of $-2.5$ to $-1\%$ in the TP2 scenario.

6 Conclusions

In this paper, we have computed the full $O(\alpha_s^2\alpha)$ corrections to gluino pair production at the LHC and at the Tevatron. Two different renormalization schemes were used. The numerical value of the $O(\alpha_s^2\alpha)$ contribution is rather independent on the renormalization scheme. The treatment of the IR and collinear singularities was performed within two different methods.

We have studied the numerical impact of the $O(\alpha_s^2\alpha)$ contributions at the LHC in two different scenarios and we have performed scans over many regions of the parameter space. The EW corrections are negative and can be safely neglected. Compared to squark–anti-squark [25,28] and squark-gluino [29] production, the EW contributions to gluino pair production are less important. The main reason is that the EW contributions do not enter the gluon fusion channel, which is the leading tree-level production channel in a wide part of the region of the parameter space investigated in this paper.

We have also provided numerical results for gluino pair production at the Tevatron selecting two scenarios belonging to the region of the parameter space investigated by the D0 and CDF collaborations. Again, the $O(\alpha_s^2\alpha)$ contributions are small and negligible.

Acknowledgments

We are indebted to Jan Germer, Wolfgang Hollik, and Maike Trenkel for useful discussions and for reading the manuscript. Michael Rauch is gratefully acknowledged for countless suggestions in the early stages of this work.
Appendix

A  Feynman Diagrams

In this appendix we collect the relevant Feynman diagrams. In the following the label $S$ ($S^0$) is used to denote charged (neutral) Higgs bosons. Moreover $V^0 = \gamma, Z$.

![Figure 1: Tree-level diagrams for the processes $q\bar{q} \rightarrow \tilde{g}\tilde{g}$ and $gg \rightarrow \tilde{g}\tilde{g}$.](image1)

![Figure 2: Tree-level diagrams for the real photon emission process $q\bar{q} \rightarrow \tilde{g}\tilde{g}\gamma$.](image2)

![Figure 3: Tree-level diagrams for the process $q\gamma \rightarrow \tilde{g}\tilde{g}q$. The diagrams for the process $\gamma\bar{q} \rightarrow \tilde{g}\tilde{g}\bar{q}$ can be obtained inverting the arrows.](image3)
Figure 4: One-loop EW diagrams for the process $q \bar{q} \rightarrow g \bar{g}$. Diagrams with crossed final states are not shown.
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Figure 5: Invariant mass distribution for gluino pair production at the LHC. In the left panels we show the LO (black line) and the LO + NLO EW (red line) distribution. The two lines are indistinguishable, owing to the smallness of the EW contributions. In the right panels the electroweak correction relative to the total result is shown.
Figure 6: Same as Fig. 5, but considering the transverse momentum distribution.
Figure 7: $\Delta = \sigma_{\text{EW, NLO}} / (\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}) \cdot 100$ as a function of $A_t$ and $A_b$ for different values of $\tan \beta$ and $M_{\text{Susy}}$. The other parameters are fixed to their SPS1a' values.

Figure 8: $\Delta = \sigma_{\text{EW, NLO}} / (\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}) \cdot 100$ as a function of $\tan \beta$ and $M_{A_0}$ for different values of $m_{\tilde{g}}$ and $M_{\text{Susy}}$. The other parameters are fixed to their SPS1a' values.
Figure 9: $\Delta = \frac{\sigma_{\text{EW, NLO}}}{\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}} \cdot 100$ as a function of $\mu$ and $M_2$ for different values of $m_{\tilde{g}}$ and $M_{\text{Susy}}$. The other parameters are fixed to their SPS1a' values.

Figure 10: $\xi \equiv -\Delta = -\frac{\sigma_{\text{EW, NLO}}}{\sigma_{\text{QCD, LO}} + \sigma_{\text{EW, NLO}}} \cdot 100$ as a function of $M_{\text{Susy}}$ and $m_{\tilde{g}}$ for different values of $M_2$ and $\tan \beta$. The other parameters are fixed to their SPS1a' values.
Figure 11: Invariant mass distribution of the two gluinos produced at the Tevatron via the process $p\bar{p} \rightarrow \tilde{g}\tilde{g}X$. In the left panels we show the LO and the EW+NLO EW contribution, while the electroweak corrections relative to the total result are shown in the right panels.
Figure 12: Same as Fig. 11, but considering the transverse momentum distribution.