Supersymmetric NLO QCD Corrections to Resonant Slepton Production and Signals at the Tevatron and the LHC

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We compute the total cross section and the transverse momentum distribution for single charged slepton and sneutrino production at hadronic colliders including NLO supersymmetric and non-supersymmetric QCD corrections. The supersymmetric QCD corrections can be substantial. We also resum the gluon transverse momentum distribution and compare our results with two Monte Carlo generators. We compute branching ratios of the supersymmetric decays of the slepton and determine event rates for the like-sign dimuon final state at the Tevatron and at the LHC.

I. INTRODUCTION

Supersymmetry [1] is the most promising and widely studied solution to the hierarchy problem [2, 3, 4, 5] of the Standard Model of particle physics (SM) [6, 7]. It must be broken with a mass scale \( O(1 - 10 \text{ TeV}) \) [8, 9, 10] and thus should be testable at the Tevatron and the LHC. When extending the SM by supersymmetry, we must introduce an additional Higgs doublet, as well as double the spectrum. The most general renormalizable superpotential consistent with this minimal particle content as well as the gauge symmetry of the SM is [11, 12]

\[
W = W_{\text{PS}} + W_{\text{HS}},
\]

\[
W_{\text{PS}} = h^{ij}_L L_i H_d E_j + h^{ij}_Q Q_i H_d D_j + h^{ij}_H Q_i H_u U_j + \mu H_d H_u,
\]

\[
W_{\text{HS}} = \lambda_{ijk} L_i L_j E_k + \lambda_{ijk} L_i Q_j D_k + \lambda_{ijk} U_i D_j H_u + \kappa_i L_i H_u.
\]

Here we use the standard notation of [13], in particular: \( i, j, k = 1, 2, 3 \) are generation indices. This superpotential leads to rapid proton decay via the \( LQD \) and \( UDD \) operators [14] in disagreement with the lower experimental bound [13]. The supersymmetric SM thus requires an additional symmetry to stabilize the proton. It was recently shown, that there are only three such discrete symmetries which are consistent with an underlying anomaly-free \( U(1) \) gauge theory [14, 15, 16]: R-parity (or equivalently matter parity), baryon-triality, and proton-hexality. R-parity prohibits \( W_{\text{PS}} \) and thus apparently stabilizes the proton. However, it allows dangerous dimension-five proton decay operators such as \( QQQL \) [17]. This long-standing problem in supersymmetry is resolved by proton-hexality, \( P_6 \), a discrete \( Z_6 \)-symmetry, which acts like R-parity on Eq. (I.1) but in addition disallows the dangerous dimension-five operators. An equally well motivated solution to the proton decay problem is baryon-triality, a discrete \( Z_3 \)-symmetry, which prohibits the \( UDD \) operator in Eq. (I.1). The baryon-triality collider phenomenology has three main distinguishing features compared to the \( P_6 \)-MSSM [13, 20] (minimal supersymmetric standard model).

1. The lightest supersymmetric particle (LSP) is no longer stable.

2. Supersymmetric particles can be produced singly, on resonance.

3. Lepton flavour and lepton number are violated.

These lead to dramatically different signatures at hadron colliders [21, 22]. It is the purpose of this paper to investigate specific baryon-triality processes at the Tevatron and the LHC in detail. We focus on the resonant production of charged sleptons and sneutrinos via the operator \( LQD \) of Eq. (I.1) and their subsequent decay into observable final states at the Tevatron and the LHC. This is the only novel production process and it has the highest kinematic reach in the supersymmetric masses.

Since the LSP is not stable, it need not be the lightest neutralino as in the MSSM, e.g. it could be a scalar tau [22]. However, the conventional supersymmetric parameter points considered in collider studies to date, the SPS points [24], imply a neutralino LSP and we restrict ourselves to this case here. We therefore consider the following reactions

\[
\tilde{d}_j d_k \rightarrow \tilde{\nu}_i \to \begin{cases} \tilde{d}_j d_k, & (a) \\ \nu_\chi^0_{-n}, & (b) \\ \ell^+ \chi^0_n, & (c) \end{cases}
\]

\[
u_j \tilde{d}_k \rightarrow \tilde{\ell}^+_i \to \begin{cases} \tilde{u}_j \tilde{d}_k, & (d) \\ \ell^- \chi^0_{n}, & (e) \\ \tilde{\nu}_i \chi^\pm_n, & (f) \end{cases}
\]
Here $\tilde{\nu}_i, \tilde{\ell}_i^\pm$ denote the sneutrino and the charged slepton, and $\tilde{\chi}^0_m, \tilde{\chi}^\pm_n$ denote the neutralino and chargino mass eigenstates, respectively. The decays (124, d) are via the $LQD$-operator. The decays (124, b, c, e, f) are cascade decays typically via a gauge coupling. In 1997 ZEUS and H1 observed a slight anomaly in their high-$Q^2$ data [25, 26], which could potentially be interpreted in terms of resonant squark production via the $LQD$ operator [27, 28, 29]. However, after computing the NLO QCD corrections to squark production at HERA [30] and also comparing with NLO QCD corrections to leptoquark pair production at the Tevatron [31], this interpretation was not consistent with the Tevatron data. We take this as a motivation to compute the NLO QCD and supersymmetric QCD corrections to the resonant slepton total and differential production cross section at the Tevatron and the LHC.

Resonant slepton production at hadron colliders via the $LQD$ operator was first investigated in [32, 33]. Using tree-level production cross sections and the decays of Eq. (1.2), a detailed phenomenological analysis of Eq. (1.2) was performed including a HERWIG [37, 38, 39, 40] simulation, again with tree-level cross sections. The NLO QCD corrections to the cross section were first computed in [41, 42]. We go beyond this work in several aspects. First, we complete the calculation by including the SUSY-QCD corrections. These involve virtual gluinos as well as tri-linear scalar couplings due to the supersymmetry breaking $A$-terms [43]. As we shall see, these can be substantial in certain regions of parameter space and modify the QCD prediction by up to 35%. We have also resummed the $p_T$ distribution of the final state gluons, c.f. [12]. In order to facilitate the implementation of our results in experimental investigations we present a detailed comparison with the HERWIG and SUSYGEN [44] Monte Carlo programs. Furthermore, we present a detailed phenomenological analysis of the final state including the relevant decay branching ratios. We also present event rates at the LHC and Tevatron for the promising like-sign dilepton signature using our improved calculation.

In related work, resonant squark production at hadron colliders via the $UDD$ operator at tree level was discussed in [33]. The NLO corrections and the $p_T$ resummation has been studied in [45]. At HERA resonant squark production has been considered in [20, 46, 47]. At $e^+e^-$ colliders one can have resonant sneutrino production via the $L_1L_i\tilde{E}_l$-operator. This has been considered in [32, 48, 49]. At hadron colliders resonant production proceeds via the operators $\lambda_{ijk}L_iQ_jD_k$. However, for $j, k = 2, 3$ the production rate is suppressed due to the lower luminosity for sea quarks. The current best bounds for the first generation quarks are given by [20, 50, 51].

| Coupling | $\lambda_{111}$ | $\lambda_{211}$ | $\lambda_{311}$ |
|----------|----------------|----------------|----------------|
| best bound | 0.0005 | 0.059 | 0.11 |

The last two bounds scale with the relevant scalar fermion mass ($m_\tilde{f}/100$ GeV). The first strict bound arises from neutrinoless double beta decay searches and has a more complicated $m_\tilde{f}$ dependence.

The outline of our paper is as follows. In Sect. II we compute the NLO resonant slepton production cross section at hadron colliders, including both virtual and real diagrams. We also employ this calculation to determine the running of the parameter $\lambda_{ijk}$. We then first compute the partonic cross section and then numerically the hadronic cross section. In Sect. III we investigate the transverse momentum distribution of the produced sleptons at NLO. In particular we perform a $p_T$ resummation of the gluon distribution and then compare with the HERWIG and SUSYGEN Monte Carlo generators. In Sect. IV we study the branching ratios of the various slepton decays and determine the resulting like-sign dimuon event rates at the LHC and the Tevatron. In Sect. V we summarize our results and conclude.

II. TOTAL PRODUCTION CROSS SECTION

A. LO Processes

We first consider the single slepton production processes at leading order, Fig. II(a),

$$d_k(p) + \tilde{d}_j/\tilde{u}_j(p') \rightarrow \tilde{\nu}_i/\tilde{\ell}_L(q),$$

$$\tilde{d}_k(p') + d_j/\tilde{u}_j(p) \rightarrow \tilde{\nu}_i/\tilde{\ell}_L(q).$$

(II.1)

The relevant parts of the Lagrangian for the production at hadronic colliders follow from the superpotential Eq. (1.1). Expressed in terms of the (four-) component fields and using the projection operators $P_R/L = \frac{1}{2}(1 \pm \gamma^5)$ we have

$$\mathcal{L}_{LQD} \supset \mathcal{X}_{ijk} (\tilde{\nu}_i \tilde{d}_k P_L d_j - \tilde{\ell}_L \tilde{u}_k P_L u_j) + \mathcal{X}_{(ijk)} (\tilde{\nu}_i \tilde{d}_j P_R d_k - \tilde{\ell}_L \tilde{u}_j P_R d_k).$$

(II.2)

We obtain for the total partonic single slepton (charged slepton or sneutrino) production cross section [32, 33]

$$\hat{\sigma}^{LO} = \frac{|\mathcal{X}_{ijk}|^2}{128} \pi (1 - \tau) \equiv \hat{\sigma}_0 (1 - \tau)$$

(II.3)

with $\tau = \tilde{m}^2/\hat{s}$, where $\tilde{m}$ is the slepton mass and $\sqrt{\hat{s}}$ the partonic center-of-mass energy.

B. NLO SUSY-QCD Corrections

The cross section at $\mathcal{O}(\alpha_s)$ comprises radiative corrections to the quark-antiquark process including virtual gluon, gluino, quark and squark exchange, Fig. II(b), (c), (d), (e), real gluon radiation,
Fig. 1(f), and Compton-like gluon initiated subprocesses, Fig. 1(g). In the following, we shall concentrate on positively charged slepton production. The production of sneutrinos can be treated analogously. Throughout this section, generation indices are suppressed.

We use dimensional regularization to isolate ultraviolet, infrared and collinear divergences, and adopt the MS scheme for the renormalization of the coupling $\lambda'$ and the factorization of the collinear divergences through a redefinition of the parton distribution functions.

The calculation of the NLO QCD corrections from virtual gluon exchange and real gluon emission is straightforward. For the partonic $q\bar{q}'$ cross section at NLO we find

$$\sigma_{q\bar{q}'}^{\text{NLO, QCD}} = \sigma_0 \delta(1-\tau) \left\{ 1 + \frac{\alpha_s}{\pi} C_F \left( \frac{3}{2} \ln \frac{\mu_R^2}{\bar{m}^2} - \frac{3}{2} \ln \frac{\mu_F^2}{\bar{m}^2} + \frac{\pi^2}{3} - 1 \right) \right\}$$

$$+ \sigma_0 \frac{\alpha_s}{\pi} C_F \left\{ \frac{1 + \tau^2}{[1 - \tau]^+} \ln \frac{\bar{m}^2}{\tau \mu_F^2} + 2(1 + \tau^2) \left[ \ln \frac{1 - \tau}{\tau} \right] \right\},$$

(II.4)

and for the quark-gluon scattering contribution

$$\sigma_{qg}^{\text{NLO}} = \sigma_0 \frac{\alpha_s}{4\pi} \left\{ \frac{1}{2} (1 - \tau)(7\tau - 3) \right\}$$

$$+ \left[ \tau^2 + (1 - \tau)^2 \right] \left[ \ln \frac{\bar{m}^2}{\mu_F^2} + \ln \frac{1 - \tau}{\tau} \right],$$

(II.5)

where $C_F = 4/3$ and $\mu = \bar{m}^2/s$. The couplings $\lambda'(\mu_R)$ and $\alpha_s(\mu_F)$ are evaluated at the renormalization scale $\mu_R$, and $\mu_F$ denotes the factorization scale. Our results agree with those presented in Ref. 44.

For a complete cross section prediction at $\mathcal{O}(\alpha_s)$ one has to include SUSY-QCD corrections through virtual gluino and squark exchange. These corrections have not been considered in previous calculations. The SUSY-QCD corrections at NLO arise from the vertex correction, Fig. 1(d), and the quark self energy, Fig. 1(e), and are infrared-finite.

The interaction of quarks $q$, (right-) left-handed squarks $\tilde{q}_L$ ($\tilde{q}_R$) and gluinos $\tilde{g}$ is described by

$$\mathcal{L}_{qq\tilde{q}}^{\text{soft}} = -\frac{1}{4} \lambda'_{ijk} A \left( \tilde{q}_{Lj} \tilde{d}_{Rk} - \tilde{q}_{Lj} \tilde{u}_{Rk} \right) \tilde{q}_{Lj} \tilde{d}_{Rk} - \frac{1}{2} \lambda'_{ijk} A \left( \tilde{q}_{Lj} \tilde{d}_{Rk} - \tilde{q}_{Lj} \tilde{u}_{Rk} \right) \tilde{d}_{Rk} \tilde{d}_{Rk}.$$

(II.7)

We have assumed here that the soft breaking terms have a universal dimensionful parameter $A$ and are proportional to the dimensionless coupling constant $\lambda'_{ijk}$ of the superpotential 54.

Due to angular momentum conservation only those diagrams contribute where the incoming quark and antiquark have the same helicity. Therefore, for a specific slepton production process, only a single vertex correction diagram contributes at NLO which is described by Eq. 17.

Using standard notation 55 for scalar loop integrals, the SUSY vertex contribution to the NLO cross section can be written as

$$\sigma_{\text{SUSY, vc}}^{\text{NLO}} =$$

$$-\frac{\alpha_s}{\pi} C_F A \tilde{m} \hat{C}_0 (-p, p', m_{\tilde{g}}, m_{\tilde{q}L}, m_{\tilde{q}R}),$$

(II.8)

where $p, p'$ are the momenta of the incoming quark and antiquark, $m_{\tilde{g}}$ is the gluino mass, and $m_{\tilde{q}L}, m_{\tilde{q}R}$ are the masses of the left- and right-handed squark and antisuark running in the loop.

1. **Vertex correction**

For the SUSY-QCD vertex corrections, also the interaction vertices of two squarks and a (left-handed) SUSY-QCD self energy contributions arise from diagrams of the form Fig. 1(e). Since the SUSY partners of left- and right-handed quarks can have different
masses, the SUSY self energy corrections to the left- and the right-handed quark field differ in general. The field strength renormalization constant $\hat{Z}$ is

$$\hat{Z}^{-1} = 1 - \left( \hat{\Sigma}_V - \hat{\Sigma}_V \big|_{UV} \right),$$

where the vector part of the self energy $\hat{\Sigma}_V$ can be written in terms of the coefficient function $B_1$ as

$$\hat{\Sigma}_V = -\frac{\alpha_s}{2\pi} C_F B_1(-p, m_{\tilde{q}}, m_{\tilde{q}L/R}).$$

The UV pole of the integral is, in the $\overline{\text{MS}}$ scheme,

$$B_1 \big|_{UV} = -(4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{2\epsilon},$$

where the number of space-time dimensions is $N = 4 - 2\epsilon$.

The SUSY-QCD self energy contributions to the NLO production cross section follow directly,

$$\hat{\sigma}_{\text{SUSY}, ss}^{\text{NLO}} = \hat{\sigma}_{\text{LO}} \frac{\alpha_s}{2\pi} C_F \times \left( B_1(-p, m_{\tilde{g}}, m_{\tilde{q}L}) + B_1(-p, m_{\tilde{g}}, m_{\tilde{q}R}) - 2B_1 \big|_{UV} \right).$$

This completes the computation of resonant slepton production at NLO.

### C. The running coupling $\lambda'$

In this section, we briefly discuss the $O(\alpha_s)$ running of the $LQD$ coupling $\lambda'$ including both QCD and SUSY-QCD corrections.

The bare $LQD$ coupling $\lambda^0$ is related to the renormalized coupling via

$$\lambda^0 = Z \lambda' \mu^\epsilon,$$

where $Z$ is the renormalization constant for the coupling $\lambda'$ and the scale $\mu$ is introduced to keep the renormalized coupling dimensionless. Note that the only scale dependence of $Z$ is carried by $\alpha_s$ and, therefore,

$$\frac{dZ}{d\mu} = \frac{dZ}{d\alpha_s} \frac{d\alpha_s}{d\mu} = -\beta(\alpha_s) \frac{dZ}{d\alpha_s},$$

where the $\beta$-function is defined as $\beta(\alpha_s) = -\mu \frac{d\alpha_s}{d\mu}$ with $\beta(\alpha_s) = 2\alpha_s + O(\alpha_s^2)$. We thus obtain for the $\beta$-function of the $LQD$ coupling $\lambda'$

$$\beta(\lambda') \equiv -\mu \frac{d\lambda'}{d\mu} = -\lambda' \beta(\alpha_s) \frac{dZ}{d\alpha_s} + \lambda' \epsilon. $$

Considering only QCD corrections, the renormalization constant $Z$ depends on the (identical) renormalization constants for the left-/right-handed quark field $Z_{L/R}$ and on the vertex renormalization $Z_{\lambda'}$. They are

$$Z_N = 1 - \frac{\alpha_s}{4\pi} C_F (4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{\epsilon},$$

$$Z_{L/R} = 1 - \frac{\alpha_s}{4\pi} C_F (4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{\epsilon}.$$  

Taking also the SUSY-QCD corrections into account, SUSY self energy contributions $\hat{Z}_{L/R}$ alter the renormalization constant $Z$,

$$\hat{Z}_{L/R} = 1 - \frac{\alpha_s}{4\pi} C_F (4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{\epsilon}.$$  

Explicitly, including only QCD corrections,

$$Z = \frac{Z_N}{Z_L} = 1 - \frac{3\alpha_s}{4\pi} C_F (4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{\epsilon},$$

and including QCD and SUSY-QCD corrections,

$$Z = \frac{Z_{N'}}{Z_L Z_L} = 1 - \frac{\alpha_s}{2\pi} C_F (4\pi)^{\epsilon} \Gamma(1 + \epsilon) \frac{1}{\epsilon}. $$
For the $\beta$-function Eq. (II.15) one obtains
\[\beta(\lambda') = (\beta^\text{QCD}_{\lambda'} + \beta^\text{SUSY}_{\lambda'}) \lambda' \frac{a_s}{2\pi} + \mathcal{O}(\lambda' a_s^2), \] (II.21)
with the coefficients $\beta^\text{QCD}_{\lambda'}, \beta^\text{SUSY}_{\lambda'}$ related to QCD and SUSY-QCD effects, respectively,
\[\beta^\text{QCD}_{\lambda'} = 3C_F, \quad \beta^\text{SUSY}_{\lambda'} = -C_F, \] (II.22)
such that in total $\beta^\text{tot}_{\lambda'} = \beta^\text{QCD}_{\lambda'} + \beta^\text{SUSY}_{\lambda'} = 2C_F$.

By integrating Eq. (II.15) and inserting the definition such that in total $\beta^\text{tot}_{\lambda'} = \beta^\text{QCD}_{\lambda'} + \beta^\text{SUSY}_{\lambda'} = 2C_F$.

D. Decoupling of heavy SUSY particles

In the limit of heavy SUSY particles $m_{\tilde{q}}$, $m_{\tilde{g}}$, $m_{\tilde{g}} \equiv \tilde{M}$ and $\tilde{M} \to \infty$, the NLO SUSY-QCD vertex corrections vanish
\[\hat{\sigma}^\text{NLO}_{\text{SUSY}, \text{vc}} = -\hat{\sigma}_{\text{LO}} \frac{\alpha_s}{\pi} C_F A \tilde{M} C_0(-p, p', \tilde{M}, \tilde{M}, \tilde{M}) \to 0, \] (II.24)
while the self energy contributions are logarithmically enhanced
\[\hat{\sigma}^\text{NLO}_{\text{SUSY}, \text{se}} = \hat{\sigma}_{\text{LO}} \frac{\alpha_s}{2\pi} C_F [B_1(-p, \tilde{M}, \tilde{M}) + B_1(-p, \tilde{M}, \tilde{M})] \to -\hat{\sigma}_{\text{LO}} \frac{\alpha_s}{2\pi} C_F \ln \frac{\mu^2}{M^2}. \] (II.25)

The large logarithms can be absorbed by a redefinition of the coupling $\lambda'$ according to
\[\lambda' = \lambda'(\mu_0) \left[1 - \frac{\alpha_s}{4\pi} \beta^\text{tot}_{\lambda'} \ln \frac{\mu^2}{\mu_0^2}\right] \] \[\downarrow \] \[\lambda' + \Delta \lambda' = \lambda'(\mu_0) \left[1 - \frac{\alpha_s}{4\pi} \beta^\text{tot}_{\lambda'} \ln \frac{\mu^2}{\mu_0^2} - \frac{\alpha_s}{4\pi} C_F \ln \frac{\mu^2}{M^2}\right]. \] (II.26)

Then, the running of $\lambda'$ is determined only by the QCD corrections (light particles)
\[\frac{\partial (\lambda' + \Delta \lambda')}{\partial \ln \mu^2} = -\lambda'(\mu_0) \frac{\alpha_s}{4\pi} \left[\beta^\text{tot}_{\lambda'} + C_F\right] \] \[= -\lambda'(\mu_0) \frac{\alpha_s}{4\pi} \beta^\text{QCD}_{\lambda'}. \] (II.27)
Expanding $\lambda' = (\lambda' + \Delta \lambda') - \Delta \lambda'$ in Eq. (II.26) cancels the term $\propto \ln(\mu^2/M^2)$ so that in the limit $\tilde{M} \to \infty$ the SUSY-QCD corrections vanish. Note that we will not adopt the decoupling scheme described above in the subsequent numerical analysis as we do not consider scenarios where squark and gluino masses are much larger than the slepton mass.

E. Total Hadronic Cross Sections

In this section, we will present numerical results for the total hadronic cross section for resonant charged slepton production at the Tevatron and at the LHC. We assume only a single non-zero coupling $\lambda'_{11} = \lambda'_{ud}$ and set this coupling $\lambda' = 0.01$ at all scales. The cross sections are directly proportional to $|\lambda|^2$ such that results for other values of the coupling are easily obtained by rescaling. For the calculation of the $PP$ and $PP$ cross sections we have adopted the CTEQ6L1 and CTEQ6M [67] parton distribution functions at LO and NLO, corresponding to $\Lambda^\text{LO} = 165$ MeV and $\Lambda^\text{NLO} = 226$ MeV at the one- and two-loop level of the strong coupling $\alpha_s(\mu_R)$, respectively. The renormalization and factorization scales have been identified with the slepton mass, $\mu_R = \mu_F = m$.

The hadronic production cross section for the process $u\bar{d} \to \tilde{\ell}^+\tilde{\ell}$ at the Tevatron and at the LHC is shown as a function of the slepton mass in Fig. 2.

We compare the LO result with the NLO results with only QCD and both QCD and SUSY-QCD corrections included. The NLO cross section for producing sleptons with a mass $m = 100$ GeV is approximately 1 pb at the Tevatron and 10 pb at the LHC, decreasing rapidly with increasing slepton mass. The slepton discovery reach at the Tevatron is thus limited by the comparably small cross section, with only one expected event at $m = 600$ GeV (for $\lambda'_{11} = 0.01$) for an integrated luminosity of $\int dt \mathcal{L} = 1$ fb$^{-1}$. The cross section is more favourable at the LHC, where a 600 GeV slepton would be produced with $\sigma \approx 50$ fb, leading to 500 events per year even for the low luminosity phase with $\mathcal{L} = 10$ fb$^{-1}$/year.

To demonstrate the impact of the SUSY-QCD cor-
FIG. 3: SUSY-QCD $K$-factor for resonant slepton production via $\lambda'_{ij1}=0$ at the LHC, $\sqrt{S}=14$ TeV, for three sets of squark and gluino masses ($m_{\tilde{q}_L}=m_{\tilde{q}_R}=m_{\tilde{g}}$). Different values for the SUSY breaking parameter $A$ are chosen as indicated. For comparison we give also the QCD $K$-factor (solid lines). The CTEQ6 LO/NLO PDFs have been used, and renormalization and factorization scales have been identified with the slepton mass $\tilde{m}$.

FIG. 4: The same as Fig. 3 but for the Tevatron.

rections more clearly, we consider the $K$-factor $K = \sigma_{\text{NLO}}/\sigma_{\text{LO}}$ defined as the ratio of NLO to LO cross section, with all quantities calculated consistently in lowest and next-to-leading order. In Fig. 3–4 we display the QCD as well as the SUSY-QCD $K$-factor at the LHC (Tevatron) for three different sets of squark and gluino masses and for three different values of the SUSY breaking parameter $A$.

The QCD $K$-factor shows an enhancement of the LO cross section of up to 35% at the LHC and of up to 50% at the Tevatron. For scenarios where the trilinear SUSY breaking parameter $A$ vanishes, the SUSY-QCD corrections are moderate. They enhance the cross section prediction by up to 10% for light sleptons and reduce it by 5 – 10% for heavier sleptons. For non-vanishing $A$ there are additional contributions due to the SUSY-QCD vertex corrections which can lead to a significant further enhancement or reduction of the cross section by up to 15% for $m_{\tilde{q},\tilde{g}} = 300$ GeV. For light squark masses, it can be as much as 35%. One observes the usual resonance peak when the slepton mass is twice the mass of the squarks in the vertex loop.

We remark that the inclusion of the NLO-QCD corrections reduces the factorization scale dependence of the theoretical prediction as expected. At the Tevatron, we consider the slepton mass interval from 100 GeV to 1 TeV. A variation of the factorization scale
between $\mu_F = 2\tilde{m}$ and $\mu_F = \tilde{m}/2$ then changes the LO cross section by up to $\pm 20\%$. This uncertainty is reduced to less than $\pm 10\%$ at NLO. At the LHC, we consider slepton masses between 100 GeV and 2 TeV and again vary the factorization scale between $\mu_F = 2\tilde{m}$ and $\mu_F = \tilde{m}/2$. The LO uncertainty is then up to $\pm 10\%$ which is reduced to less than $\pm 5\%$ in NLO.

We conclude this section by tabulating some selected cross section predictions for slepton production with $\tilde{m} = 300$ GeV at the Tevatron or the LHC via the $u\bar{d}$ initial state in Table I. We show results for the default CTEQ6M parton density functions (PDFs) and for the MRST 58, 59 parametrization. We find that the difference between the two PDF parametrizations is less than approximately 5% for the NLO cross section predictions.

III. SLEPTON TRANSVERSE MOMENTUM DISTRIBUTION

In this section we compute the transverse momentum ($p_T$) distribution of the produced sleptons. At Born level, momentum conservation enforces the produced sleptons to have zero $p_T$. At $O(\alpha_s)$ the sleptons may have non-zero $p_T$ due to an additional parton in the final state. The differential hadronic cross section in $p_T$ at $O(\alpha_s)$ is given in Appendix A. However, for $p_T \ll \tilde{m}$, with $\tilde{m}$ the slepton mass, the emitted parton becomes soft and/or collinear to one of the incoming partons leading to large logarithms in the differential cross section. The perturbation series in $\alpha_s$ breaks down and has to be replaced by a series in $\alpha_s^0 \ln^m(\tilde{m}^2/p_T^2)$ with $m,n \in \mathbb{N}$ and $m = 0, \ldots, 2n-1$. We shall employ the Collins, Soper, and Sterman (CSS) formalism to resum the large logarithms 60, 61, 62.

The $p_T$ spectrum of the sleptons has also been studied in 42. We present the main formulae and extend the studies of 42 by a comparison with Monte Carlo simulations.

A. Soft and Collinear Gluon Summation

The summed differential cross section, valid in the region of soft and/or collinear emitted partons, is given by 62

$$
\frac{d^2\sigma_{\text{resum}}}{dp_T dy} = \frac{\hat{\sigma}_0 p_T n^2}{S} \int_0^\infty db \ 0(b p_T) W(b), \quad (\text{III.1})
$$

where $S$ is the hadronic c.m. energy, $J_0$ the zeroth order Bessel function, and $\hat{\sigma}_0 = \hat{\sigma}_0(s = \tilde{m}^2)$ is the leading order cross section defined in Eq. (II.3). The Sudakov-like form factor $W(b)$ is given by

$$
W(b) = \exp \left\{ -\int_{(b_0/b)^2}^{\tilde{m}^2} \frac{dq^2}{q^2} \frac{\alpha_s(q^2)}{2\pi} \left[ K_1 \ln \frac{\tilde{m}^2}{q^2} + K_2 \right] \right\} 
\times \left[ f_{qA}(x^0_1) f_{qB}(x^0_2) + (x^0_1 \leftrightarrow x^0_2) \right], \quad (\text{III.2})
$$

where $x^0_{1/2} = e^{\pm y} \sqrt{\tilde{m}^2/S}$ are the parton momentum fractions in the limit $p_T/\tilde{m} \to 0$ and $y$ denotes the slepton rapidity. The $f_{qA}$, $f_{qB}$ are the parton density functions for hadrons $A$, $B$, respectively. The coefficients $K_1$, $K_2$ are in general functions of $\alpha_s$. However, as we discuss below, we are here only interested in the exponential at $O(\alpha_s)$, i.e. $K_{1,2}$ are constant. The integration in Eq. (III.2) has to be performed over the impact parameter $b$, the Fourier conjugate of $p_T$. We perform the integration analytically by using the LO running of $\alpha_s$. The PDF’s are evaluated at a fixed factorization scale $\mu_F$, where the most convenient choice 62 is $\mu_F = b_0/b$ with $b_0 = 2e^{-\gamma_E}$ and $\gamma_E$ the Euler-Mascheroni constant.

We calculate the coefficients $K_{1,2}$ by expanding (III.1) up to $O(\alpha_s)$ and also for some algebra one obtains 63.

$$
\frac{d^2\sigma_{\text{asym}}}{dp_T dy} = \frac{\hat{\sigma}_0 \tilde{m}^2}{S p_T} \left\{ \left[ K_1 \ln \frac{\tilde{m}^2}{p_T^2} + K_2 \right] f_{qA}(x^0_1) f_{qB}(x^0_2) \right. 
+ \sum_{i=0,g} f_{qA}(x^0_1) (P_{q,i} \circ f_{iB})(x^0_2) + \sum_{i=0,g} f_{qA}(x^0_1) (P_{qi} \circ f_{iB})(x^0_2) \right\} + (x^0_1 \leftrightarrow x^0_2), \quad (\text{III.3})
$$

where $P_{q,i}$, $P_{qi}$ are the Altarelli-Parisi splitting functions 44, and the convolution is defined as

$$
(f \circ g)(x) = \int x \frac{dz}{z} f(z) g\left(\frac{x}{z}\right). \quad (\text{III.4})
$$

We can then compare the soft- and collinear divergences of this asymptotic cross section with those of the perturbative result in the limit $p_T \to 0$. We obtain $K_1 = 2C_F$ and $K_2 = -3C_F$ with $C_F = 4/3$.

It was pointed out in 63 that $W(b)$ is ill-defined when $b > 1/\Lambda_{QCD}$, because confinement sets in. We factor out the non-perturbative part by replacing $W(b)$ by
The full transverse momentum distribution for the case of a (left-handed) slepton at the Tevatron (LHC) is shown in Fig. 5 (Fig. 6). The $y$-dependence has been integrated out and the renormalization and factorization scales of the asymptotic and perturbative cross section have been set equal to $\mu = \sqrt{p_T^2 + \vec{m}^2}/2$. Again we set $\lambda_{i11} = 0.01$ and use the CTEQ6M PDFs [57]. We normalize the integral of the transverse momentum distribution to the total hadronic cross section including QCD corrections. In the low-$p_T$ region, the perturbative and the asymptotic cross section cancel each other and the spectrum is described well by the summed prediction. At high $p_T$, i.e. $p_T > 50$ GeV (100 GeV) at the Tevatron (LHC), the summed and asymptotic contribution start to fall off and the distribution is to a good approximation given by the perturbative cross section. The summed cross section peaks at $p_T = 3.1$ GeV ($p_T = 3.4$ GeV) at the Tevatron and at $p_T = 4.3$ GeV ($p_T = 5.0$ GeV) at the LHC for $\vec{m} = 200$ GeV ($\vec{m} = 500$ GeV).

Yanget al. [12] derived the transverse momentum distribution for sneutrinos with masses of 200 GeV, 400 GeV and 600 GeV. They did not use a matching procedure [11] and did not normalize the integral of the transverse momentum distribution to the total hadronic cross section including QCD corrections. To facilitate the comparison with [12] we have also calculated the $p_T$ distribution of the sneutrinos without a matching function and without normalizing to the total NLO cross section. In the phenomenologically relevant region, i.e. where $p_T \lesssim 50$ GeV, there is good agreement with [12]. At high $p_T$ we find discrepancies. Our differential cross section at $p_T = 100$ GeV ($p_T = 200$ GeV) is 75% (30%) larger (smaller) compared to [12] for a sneutrino with a mass of 200 GeV at the Tevatron (LHC). These discrepancies are not particularly relevant phenomenologically, because the differential cross section is between two and three orders of magnitude smaller at $p_T = 100$.

### C. Results

| TABLE I: Numerical results for the production cross section for a slepton of mass $\tilde{m} = 300$ GeV at the Tevatron ($\sqrt{s} = 1.96$ TeV) and the LHC ($\sqrt{s} = 14$ TeV). Slepton production via the $u\bar{d}$ initial state is considered, with $\lambda_{i11} = 0.01$. All scales are fixed to the central scale $\mu = \tilde{m}$. Squark and gluino masses are chosen as $m_{\tilde{q}_{i}} = m_{\tilde{g}} = m_{3} = 600$ GeV. Results are presented for the CTEQ6L1/CTEQ6M [57] and the MRST2002LO/MRST2004NLO [53,59] PDF parametrizations. |
|----------------------------------|----------|----------|----------|----------|----------|----------|----------|----------|
|                                 | LO       | QCD      | SUSY-QCD | LO       | QCD      | SUSY-QCD |
|                                 | $A = 0$ TeV | $A = 1$ TeV | $A = 0$ TeV | $A = 1$ TeV |
| $\sigma$ [fb] | $25.72$ | $38.43$ | $39.23$ | $40.22$ | $25.55$ | $37.73$ | $38.52$ | $39.50$ |
| Tevatron                         |          |          |          |          |          |          |          |          |
| $\sigma$ [fb] | $358.7$  | $467.0$  | $487.0$  | $491.6$  | $347.4$  | $471.2$  | $482.4$  | $496.2$  |
| LHC                             |          |          |          |          |          |          |          |          |

$W(b) \rightarrow W(b_{s}) e^{-S_{np}(b)}$, (III.5)

$$b_{s} \equiv \frac{b}{\sqrt{1 + b^{2}/b_{\text{max}}^{2}}}$$ (III.6)

which smoothly cuts-off the region $b > b_{\text{max}}$, and our choice of $b_{\text{max}}$ is given below. $S_{np}(b)$ parameterizes the non-perturbative part and has to be determined by experiment. We follow the approach of [67], where $e^{-S_{np}(b)}$ is approximated by a Gaussian with adequate parameters fitted to Drell-Yan data resulting in

$$S_{np}(b) = b_{s}^{2} \left[ g_{1} + g_{2} \ln \frac{m}{2Q_{0}} + g_{1} g_{3} \ln (100 x_{1}^{4} x_{2}^{4}) \right].$$ (III.7)

With $b_{\text{max}} = 0.5$ GeV$^{-1}$ and $Q_{0} = 1.6$ GeV, [67] finds $g_{1} = 0.21^{+0.01}_{-0.001}$ GeV$^{2}$, $g_{2} = 0.68^{+0.01}_{-0.002}$ GeV$^{2}$, $g_{3} = -0.60^{+0.05}_{-0.04}$. Yang et al. [12] used the approach of [68] to factor out nonperturbative physics.

### B. Matching

At low $p_T$, the summed cross section describes the spectrum accurately compared to the perturbative cross section at $O(\alpha_{s})$, Eq. (A.2). All terms singular in $p_T$ are included. For large $p_T$, also the non-singular terms must be considered. They correspond to the difference of the perturbative and the asymptotic cross section at $O(\alpha_{s})$. For large $p_T$ the summed and asymptotic cross sections are negative, which is unphysical. Therefore, we suppress their contributions with an empirical matching function [68, 70]:

$$\begin{align*}
\frac{d^{2}\sigma_{\text{full}}}{dp_{T}dy} &= \frac{1}{1 + (p_{T}/p_{T}^{\text{match}})^{4}} \left[ \frac{d^{2}\sigma_{\text{resum}}}{dp_{T}dy} - \frac{d^{2}\sigma_{\text{asym}}}{dp_{T}dy} \right] \\
&+ \frac{d^{2}\sigma_{\text{pert}}}{dp_{T}dy}. 
\end{align*}$$ (III.8)

We choose $p_{T}^{\text{match}} = \tilde{m}/6$ at the Tevatron and $p_{T}^{\text{match}} = \tilde{m}/3$ at the LHC.
FIG. 5: Differential transverse momentum distribution of the slepton at the Tevatron for the $u\bar{d}$ initiated process via $\lambda'_{111}$. The two sets of curves correspond to slepton masses of 200 and 500 GeV. The different lines of each set denote the perturbative (dot dashed), asymptotic (dashed), summed (dotted) and full (solid) differential cross section.

The Monte Carlo programs generate the $p_T$ distribution from the LO process through the parton shower, for which we use the default settings. In Fig. 7 (Fig. 8) we compare the transverse momentum distribution as predicted by the analytic resummation with HERWIG and SUSYGEN. All curves are normalized to the NLO QCD cross section. For both the Tevatron and the LHC, the spectrum generated with SUSYGEN (through the PYTHIA parton shower) is significantly softer than that of HERWIG. The analytic summation lies in between the two Monte Carlo programs. The discrepancy between the HERWIG and PYTHIA parton showers is a known problem, which also occurs, for example, for resonant single Z-boson production \cite{73}. Our analytic results might help to improve future versions of Monte Carlo generators for resonant scalar particle production.

**D. Comparison with Monte Carlo Generators**

Next, we compare our analytic results for the slepton transverse momentum distribution with the distributions predicted by two Monte Carlo event generators: HERWIG 6.510 \cite{37,38,39,40} and SUSYGEN 3 \cite{44}, which uses the PYTHIA 6.205 \cite{71,72} parton shower.

In order to determine the final state particles observed in a collider experiment as well as their event rates, we must discuss the possible decays of the resonantly produced (left-handed) sleptons. Since all SUSY particles, even the LSP, are unstable in $P_6$-violating SUSY models, their decay signatures are crucial for experimental analyses. Here we shall assume a neutralino LSP. Note that in $P_6$-violating SUSY models one can also have a scalar tau LSP which dramatically alters
the potential signatures \cite{23}.

Apart from the inverted production process, the sleptons can decay through gauge interactions. Neglecting mixing between left- and right-handed sleptons the dominant decays are, cf. Eq. (IV.2):

\[
d_j \tilde{d}_k \rightarrow \tilde{\nu}_\tau \tilde{\nu}_e \\
u_j \tilde{d}_k \rightarrow \tilde{\tau}_+ \\
u_j \tilde{\nu}_e \tilde{d}_k \\
\]

In order to quantitatively discuss the possible signal rates, we consider the parameters of the six mSUGRA SPS points \cite{24}. We calculate the spectra using the SPheno 2.2.3 \cite{74} online calculator \cite{73,75,76} with the SM input \cite{77}: \(m_t = 174.2 \text{ GeV}, m_b(m_b)_{\text{MS}} = 4.2 \text{ GeV and } \alpha_{\text{em}}(M_Z) = 0.1176\). The relevant masses of the SUSY particles are shown in Table II in GeV.

In these scenarios the lightest neutralino, \(\tilde{\chi}_1^0\), is the LSP. Possible decay modes via the \(\lambda_{ijk}L_iQ_j\tilde{d}_k\) interaction are \(46, 73, 80, 81\):

\[
\tilde{\chi}_1^0 \rightarrow \tilde{\tau}_+ \tilde{\nu}_e d_k, \\
\tilde{\tau}_+ \tilde{\nu}_e \tilde{d}_k, \\
\tilde{\tau}_+ \nu_e \tilde{d}_k. \\
\]

The heavier neutralinos \(\tilde{\chi}_2^{0,\pm}\) and the charginos \(\tilde{\chi}_1^\pm\) may also decay via \(\lambda_{ijk}'\) but these channels are strongly suppressed compared to the gauge coupling decay modes into the LSP. For example, the branching ratio (BR) for the decay \(\tilde{\chi}_2^0 \rightarrow \mu^+ \mu^- d\) with \(\lambda_{211} = 0.01\) is less than \(2 \times 10^{-6}\) in the five SPS scenarios, where the slepton can decay into an on-shell \(\tilde{\chi}_1^0\).

From the multiple decay modes we concentrate on those which lead to final states with two like-sign leptons; a signature with low background \cite{37,82,83,84,85}. Here we assume \(\lambda_{211}' \neq 0\), since dimuon final states are of great interest for experimental analyses \cite{86,87,88,89}. For our numerical studies, we neglect decay modes with branching ratios \(BR < 0.5\%

### Table II: Relevant masses of the mSUGRA spectrum for the six Snowmass points, i.e. those which are kinematically accessible in the slepton decays. All masses are in GeV, rounded to the nearest whole number.

| Points | \(\tilde{\mu}_L\) | \(\tilde{\nu}_\mu\) | \(\tilde{\tau}_1^+\) | \(\tilde{\tau}_2^+\) | \(\tilde{\tau}_3^+\) | \(\tilde{\chi}_1^0\) | \(\tilde{\chi}_2^0\) | \(\tilde{\chi}_3^0\) | \(\tilde{\chi}_1^+\) | \(\tilde{\chi}_2^+\) | \(b^0\) |
|--------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| 1a     | 202             | 186             | 134             | 206             | 185             | 97              | 181             | -               | -               | 180             | -               |
| 1b     | 339             | 329             | 197             | 345             | 318             | 162             | 307             | -               | -               | 307             | -               |
| 2      | 1459            | 1456            | 1440            | 1453            | 1450            | 123             | 235             | 478             | 492             | 235             | 493             |
| 3      | 288             | 277             | 173             | 290             | 276             | 161             | -               | -               | -               | -               | 114             |
| 4      | 449             | 441             | 260             | 415             | 388             | 121             | 226             | 401             | 416             | 226             | 418             |
| 5      | 257             | 245             | 182             | 258             | 243             | 120             | 230             | -               | -               | 230             | -               |

FIG. 7: Transverse momentum distribution of the slepton at the Tevatron generated with HERWIG (solid histogram) and SUSYGEN (dashed histogram) compared with the SM input \cite{78}. The curves are normalized to the NLO QCD cross section.

FIG. 8: Same as Fig. 7 but for the LHC.
and where there are less than 0.1 events for the given luminosity. The relevant decay chains leading to a like-sign muon pair are shown in Table III. There are two distinct baryon triality scenarios. If $\lambda'_{211} = 0$, the $\tilde{\mu}$ and $\tilde{\nu}_\mu$ will have a significant decay width through the $L_2Q_1\bar{D}_1$ operator. In contrast, if $\lambda'_{211}$ is much smaller than the gauge couplings, the $\tilde{\mu}$ and $\tilde{\nu}_\mu$ decay through $P_6$-conserving channels. For example, the $BR(\tilde{\mu}^+ \rightarrow \lambda' \rightarrow u d)$ is less than 1.1% for $\lambda'_{211} = 0.01$ in all six scenarios.

Baryon triality then manifests itself mainly through the LSP decays. Note that the $BR$s of the LSP at tree level are independent of $\lambda'_{211}$. For the event rates, we choose $\lambda'_{211} = 0.05$ and $\lambda'_{211} = 0.01$ in the following. In Table IV, we give the expected number of like-sign dimuon events ($\mu^+\mu^\pm$) for the LHC (Tevatron) with an integrated luminosity of 10 fb$^{-1}$ (1 fb$^{-1}$). We use the NLO cross sections including QCD and SUSY-QCD corrections with the triple-SUSY breaking parameter $A$ set to zero. Apart from the $\tilde{\mu}$, $\tilde{\nu}_\mu$, and LSP decays, we use the $BR$s obtained by SPHENO 2.2.3 [74].

If decays into the heavier neutralinos $\chi^0_{3,4}$ and chargino $\chi_2^+$ are kinematically forbidden (SPS1a, 1b, 3, 5), the $\tilde{\mu}$ decays dominantly into two like-sign muons associated with two jets; contributions from the $\tilde{\nu}_\mu$ decay are negligible. Here all final state particles can be detected and the slepton and gaugino masses can be reconstructed. If the $\tilde{\nu}_\mu$ and $\tilde{\mu}$ are heavier than all neutralinos and charginos (SPS2, 4), then the $\tilde{\nu}_\mu$ decay gives the main contribution and also decays with an on-shell $Z^0$ in the final state are kinematically accessible.

In the case of $\tilde{c}$ and $\tilde{t}$ production, i.e. $\lambda'_{111} \neq 0$, one obtains the same number of events for a given value of $\lambda'$, however the bounds on $\lambda'_{111}$ are much stronger [20]. If $\lambda'_{211} \neq 0$ there are more possible decay modes, because left-right mixing in the stau sector can not be neglected. Both mass eigenstates $\tilde{\tau}_1$ and $\tilde{\tau}_2$ can be produced via the $L_3Q_1\bar{D}_1$ operator and the heavier mass eigenstate can decay into the lighter one via $\tilde{\tau}_2 \rightarrow \tilde{\tau}_1Z^0$ and $\tilde{\tau}_2 \rightarrow \tilde{\tau}_1h^0$, if kinematically allowed. This results in more decay chains compared to those for $\tilde{\mu}$ and $\tilde{\nu}_\mu$. In addition, the charge of the $\tau$-jets has to be reconstructed.

At the LHC, the beam energy will be high enough to produce single sleptons at a high rate and an excess of like-sign lepton pairs would be a hint of baryon-triality supersymmetry. Since the event rates scale with $|\lambda'_{ijk}|^2$ for $\lambda'_{ijk}$ much smaller than the gauge couplings, we can roughly estimate that the magnitude of $\lambda'_{211}$ can be tested down to $10^{-2} - 10^{-4}$, depending on the scenario.

## V. CONCLUSION

We have argued that in preparation for the LHC, the baryon-triality interactions must be considered in detail. We have computed the full NLO corrections to resonant slepton production at hadron colliders. We have found significant changes in the total and differential cross sections. In particular, we have computed the SUSY-QCD corrections for the first time and found that the effects can be large in specific regions of parameter space, modifying the QCD prediction by up to 35%. The factorization scale dependence of the cross section at the LHC (Tevatron) is also significantly reduced to about ±5% (±10%) at NLO. We have used our computation to determine the running of the $X'$-coupling without and with the decoupling of heavy SUSY particles.

Next, we considered the slepton traverse momentum distribution and resummed the final state gluons. Including these results we have performed a detailed comparison with the HERWIG and SUSYGEN/PyTHIA Monte Carlo generators. Significant discrepancies remain, similar to the case of SM $Z$-boson production. This needs to be remedied. Finally including the slepton decay branching ratios we determine the like-sign dimuon event rates at the LHC and Tevatron for various SUSY spectra. This should enable a detailed investigation by the experimental groups.

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**TABLE III: Relevant decays of resonantly produced single smuons $\tilde{\mu}$ and smuon sneutrinos $\tilde{\nu}_\mu$ into like-sign muon pairs of positive charge.**

1. $\tilde{\mu}^+ \rightarrow \mu^+ \chi_{1,2}^0$
   \[ \chi_{1,2}^0 \rightarrow \mu^+ \bar{u} d \]
2. $\tilde{\mu}^+ \rightarrow \mu^+ \chi_{2,3,4}^0$
   \[ \chi_{2,3,4}^0 \rightarrow Z^0 \chi_{1}^0 \]
   \[ \chi_{1}^0 \rightarrow \mu^+ \bar{u} d \]
3. $\tilde{\nu}_\mu \rightarrow \bar{\nu}_\mu \chi_{1,2}^\pm$
   \[ \chi_{1,2}^\pm \rightarrow W^\pm \chi_{1}^0 \]
   \[ \chi_{1}^0 \rightarrow (\nu_\mu \mu^+) (\mu^+ \bar{u} d) \]
4. $\tilde{\nu}_\mu \rightarrow \mu^+ \chi_{1,2}^0$
   \[ \chi_{1,2}^0 \rightarrow W^- \chi_{1}^0 \]
   \[ \chi_{1}^0 \rightarrow \mu^+ \bar{u} d \]
\[
X_{311} = 0.05
\]

| \(X\) | \(1a\) | \(1b\) | 2 | 3 | 4 | 5 |
|------|------|------|---|---|---|---|
| \(\bar{u}d\) | 50246 (647) | 15957 (148) | 11.0 (-) | 34972 (373) | 1146 (9.0) | 32126 (370) |
| \(\bar{\mu}B\) | 2997 (38.6) | - (-) | - (-) | - (-) | - (-) | - (-) |
| \(Z^0 \bar{u}d\) | - (-) | - (-) | 39.5 (-) | - (-) | 3289 (20.8) | - (-) |
| \(\bar{\nu}_e \nu_e \bar{u}d\) | - (-) | - (-) | 7.6 (-) | - (-) | 613 (3.9) | - (-) |
| \(\bar{W}^- \bar{u}d\) | - (-) | - (-) | 52.1 (-) | - (-) | 5060 (17.8) | 1752 (13.8) |

\[
X_{311} = 0.01
\]

| \(X\) | \(1a\) | \(1b\) | 2 | 3 | 4 | 5 |
|------|------|------|---|---|---|---|
| \(\bar{u}d\) | 2279 (29.3) | 723 (6.7) | 0.4 (-) | 1765 (18.8) | 57.8 (0.4) | 1442 (16.6) |
| \(\bar{\mu}B\) | 136 (1.7) | - (-) | - (-) | - (-) | - (-) | - (-) |
| \(Z^0 \bar{u}d\) | - (-) | - (-) | 1.6 (-) | - (-) | 134 (0.8) | - (-) |
| \(\bar{\nu}_e \nu_e \bar{u}d\) | - (-) | - (-) | 0.3 (-) | - (-) | 25.0 (0.2) | - (-) |
| \(\bar{W}^- \bar{u}d\) | - (-) | - (-) | 2.1 (-) | - (-) | 206 (0.7) | 80.1 (0.6) |

TABLE IV: Number of like-sign dimuon events, \(\mu^+\mu^- + X\), from cascade decays of \(\bar{\mu}\) and \(\bar{\nu}_e\) at the LHC (Tevatron) with an integrated luminosity of 10 fb\(^{-1}\) (1 fb\(^{-1}\)) for the six mSUGRA SPS scenarios. The expected number of events include also the charge conjugated final states. Events which are kinematically forbidden or with BR < 0.5% or where there are less than 0.1 events are denoted by a dash.

APPENDIX A: HIGH-\(p_T\) SLEPTON DISTRIBUTION AT \(O(\alpha_s)\)

The transverse momentum spectrum in the high \(p_T\) region, i.e. \(p_T > \tilde{m}\), can be calculated perturbatively. According to \(^7\)0, \(^9\)0, the perturbative hadron level cross section for the parton level process

\[
a(p_1) + b(p_2) \rightarrow \bar{L}(q) + X,
\]

can be written as:

\[
\frac{d\sigma_{\text{pert}}}{dp_T dy} = \frac{2p_T}{S} \left[ \int_{x_1}^1 \frac{dx_1}{x_1 x_2} f_{aA}(x_1) f_{bB}(x_2) \hat{s}^{ab} \hat{d}_T dt \right] (x_2 = x_2^*),
\]

\[
+ \int_{x_2}^1 \frac{dx_2}{x_2 x_1} f_{aA}(x_1^*) f_{bB}(x_2) \hat{d}^{ab} \hat{d}_T dt \right] (x_1 = x_1^*),
\]

(A.2)

Here, parton \(a\) with momentum \(p_1 = x_1 P_1\), and parton \(b\) with momentum \(p_2 = x_2 P_2\) produce a slepton (sneutrino) \((\bar{L})\) with momentum \(q\) and mass \(\tilde{m}\) and an additional particle, \(X\). Parton \(a\) originates from hadron \(A\) with momentum \(P_1\) and parton \(b\) originates from hadron \(B\) with momentum \(P_2\). \(f_{aA}(x_1)\) denotes the parton density function for the parton \(a\) in hadron \(A\) with momentum fraction \(x_1\). The other hadron level variables as functions of \(p_T\) are

\[
S = (P_1 + P_2)^2,
\]

\[
\sqrt{\tau^+} = \sqrt{p_T^2 / S + (\tilde{m}^2 + p_T^2) / S},
\]

\[
x^*_A / B = e^{\pm y} \sqrt{\tau^+}, \quad x_{1/2}^0 = e^{\pm y} \sqrt{\tilde{m}^2 / S},
\]

\[
x_{1/2}^+ = e^{\pm y} \sqrt{(\tilde{m}^2 + p_T^2) / S},
\]

\[
x_{1/2}^* = (x_{1/2}^+ x_{1/2}^- - \tilde{m}^2) / (x_{1/2}^- + x_{1/2}^+). \quad (A.3)
\]

Here \(y\) is the slepton rapidity and \(x_{1/2}^0\) are the parton momentum fractions in the limit \(p_T / \tilde{m} \rightarrow 0\). The partonic variables are given by

\[
\hat{s} = x_1 x_2 S,
\]

\[
\hat{t} = \tilde{m}^2 \left( 1 - x_{1/2}^0 \sqrt{1 + \frac{p_T^2}{\tilde{m}^2}} \right),
\]

\[
\hat{u} = \tilde{m}^2 \left( 1 - x_{1/2}^+ \sqrt{1 + \frac{p_T^2}{\tilde{m}^2}} \right). \quad (A.4)
\]

The transverse momentum spectrum of the slepton (sneutrino) is at \(O(\alpha_s)\) due to the gluon radiation and Compton-like processes. For gluon radiation, the relevant parton level processes are

\[
q(p_1) + \bar{q'}(p_2) \rightarrow \bar{L}(q) + g,
\]

\[
\bar{q}(p_1) + q(p_2) \rightarrow \bar{L}(q) + g, \quad (A.5)
\]

where \(q\) and \(q'\) are not necessarily the same flavour. The resulting differential partonic cross section is

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given by
\[
\frac{\hat{s} d\hat{s}^{q\bar{q}}}{dt} = \frac{\hat{s} d\hat{s}^{q\bar{q}}}{dt} = \frac{\alpha_s \lambda^2}{18 \hat{s}} \left[ \frac{\hat{t}}{\hat{s}} + \frac{\hat{t}}{\hat{u}} + 2 \left( \frac{m^2 \hat{t}}{\hat{u} \hat{t}} + 1 \right) \right],
\]
(A.6)
and \( \hat{s}, \hat{t}, \hat{u} \) are the parton level Mandelstam variables \([\Delta.2]\). The parton level Compton-like processes with \( a = q, \bar{q} \), \( b = q', \bar{q}' \) and \( g \) denoting a gluon are:
\[
\begin{align*}
a(p_1) + g(p_2) & \to \bar{L}(q) + b, \\
g(p_1) + a(p_2) & \to \bar{L}(q) + b.
\end{align*}
\]  
(A.7)
We obtain the differential partonic cross sections:
\[
\frac{\hat{s} d\hat{s}^{q\bar{a}}}{dt} = -\frac{\alpha_s \lambda^2}{48 \hat{s}} \left[ \frac{\hat{t}}{\hat{s}} + \frac{\hat{t}}{\hat{u}} + 2 \left( \frac{m^2 \hat{t}}{\hat{u} \hat{t}} + 1 \right) \right],
\]
(A.8)
These can then be inserted in Eq. \([A.2]\) together with the parton density functions to obtain the corresponding hadronic differential cross sections.

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