Stop Lepton Associated Production at Hadron Colliders

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Abstract

At hadron colliders, the search for R-parity violating supersymmetry can probe scalar masses beyond what is covered by pair production processes. We evaluate the next-to-leading order SUSY-QCD corrections to the associated stop or sbottom production with a lepton through R-parity violating interactions. We show that higher order corrections render the theoretical predictions more stable with respect to variations of the renormalization and factorization scales and that the total cross section is enhanced by a factor up to 70% at the Tevatron and 50% at the LHC. We investigate in detail how the heavy supersymmetric states decouple from the next-to-leading order process, which gives rise to a theory with an additional scalar leptoquark. In this scenario the inclusion of higher order QCD corrections increases the Tevatron reach on leptoquark masses by up to 40 GeV and the LHC reach by up to 200 GeV.

I. INTRODUCTION

Supersymmetry is one of the most promising candidates for physics beyond the Standard Model, and the search for supersymmetric particles is one of the most prominent tasks for current and future colliders. Usually, searches for supersymmetric particles at colliders assume the conservation of R parity. However, exact R parity is not in any way inherent to supersymmetric models, neither gauge invariance nor supersymmetry actually require it [1,2]. R parity is imposed to bypass problems with flavor-changing neutral currents, proton decay, atomic parity violation, and other experimental constraints. It is also crucial for supersymmetric cold dark matter. R-parity conservation has a huge impact on searches for supersymmetric particles at colliders: superpartners can only be produced in pairs and the final state has to include two stable LSPs. In the absence of R-parity conservation, single supersymmetric particles can be produced, which can extend the reach of colliders [3–5]. It is interesting to notice that pair production of scalar top quarks

(or sbottoms or scalar leptoquarks) in hadronic colliders is essentially model independent since the squark–gluon interaction is fixed by $SU(3)$ gauge invariance [6]. In contrast, single production takes place via an unknown R-parity violating Yukawa interaction λ. Nevertheless, the available phase space for single stop (sbottom/leptoquark) production is larger than the one for pair production, allowing the search to extend considerably the bounds on these particles [7]. Moreover, the single production can also give information on the Yukawa couplings λ.

In this letter, we assume that the superpotential exhibits a term $\lambda'_{ijk}\epsilon_{ab}L_i^a Q_j^b D_k^c$ where L (Q) stands for the lepton (quark) doublet superfield and D^c is the charge conjugates right handed quark superfield. This way scalar tops ($\tilde{t}_1$) and sbottoms ($\tilde{b}_1$) couple to quark–lepton pairs just like a scalar leptoquark. The production of single stops (sbottoms/leptoquarks) in association with a charged or neutral lepton proceeds via [8]:

$$p\bar{p}\,(pp) \to qg \to \ell\tilde{t}_1 \tag{1}$$

We study this process taking into account the SUSY-QCD next-to-leading order corrections. We show how higher order corrections not only enhance the total cross section, but also render the theoretical predictions more stable with respect to variations of the renormalization and factorization scales.

All our results, of course, apply to models containing a single leptoquark in addition to the Standard Model particles in the limit in which we decouple all supersymmetric states but the lightest stop or sbottom. This way, we can obtain the QCD corrections to the single leptoquark production and analyze their impact on the attainable Tevatron and LHC limits. In Section III we describe in detail the decoupling properties of the heavy supersymmetric states including next-to-leading order effects in linking supersymmetric with leptoquark-type observables.

<u>Conventions:</u> Throughout this paper we show consistent leading order or next-to-leading order cross section predictions, including the respective one loop or two loop strong coupling constant and the corresponding CTEQ5L or CTEQ5M1 parton densities [9]. We usually assume a scalar top (bottom/leptoquark) mass of 200 GeV, a massless lepton, and we set the R-parity violating coupling λ'_{SUSY} to unity. Furthermore, we assume that the squark couple to a quark and a lepton without any additional suppression from the squark mixing angle. Since the R-parity violating coupling as well as the mixing angle dependence are universal factors, as far as QCD corrections are concerned, they can be trivially added. If not stated otherwise, we scale the supersymmetric mass spectrum as $m_{\tilde{g}} = m_{\tilde{t}_1} + 100$ GeV, $m_{\tilde{q}} = m_{\tilde{t}_1} + 200$ GeV, and $m_{\tilde{t}_2} = m_{\tilde{t}_1} + 300$ GeV[1].

II. NEXT-TO-LEADING ORDER CROSS SECTION

The leading order partonic cross section for the single stop (sbottom/leptoquark) production is given by:

$$\frac{d\hat{\sigma}}{d\hat{t}} = \lambda'^2_{\mathrm{SUSY}}\frac{\alpha_s}{24\hat{s}^2}\left[-\frac{\hat{t}}{2\hat{s}} - \frac{\hat{t}\hat{u}^2}{\hat{s}(\hat{u} - m_{\tilde{t}_1}^2)^2} + \frac{\hat{u}\hat{t}}{\hat{s}(\hat{u} - m_{\tilde{t}_1}^2)}\right] \tag{2}$$

Here $\hat{s}$ is the parton–parton center-of-mass energy, $\hat{t} = (p_q - p_{\tilde{t}_1})^2$, and $\hat{u} = (p_q - p_\ell)^2$. From the definition of λ'_{SUSY} in the Lagrangean we see that left handed stops are produced only in association with charged

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[1]The FORTRAN code used in this letter is available from the authors: aalves@ift.unesp.br, tilman.plehn@cern.ch.
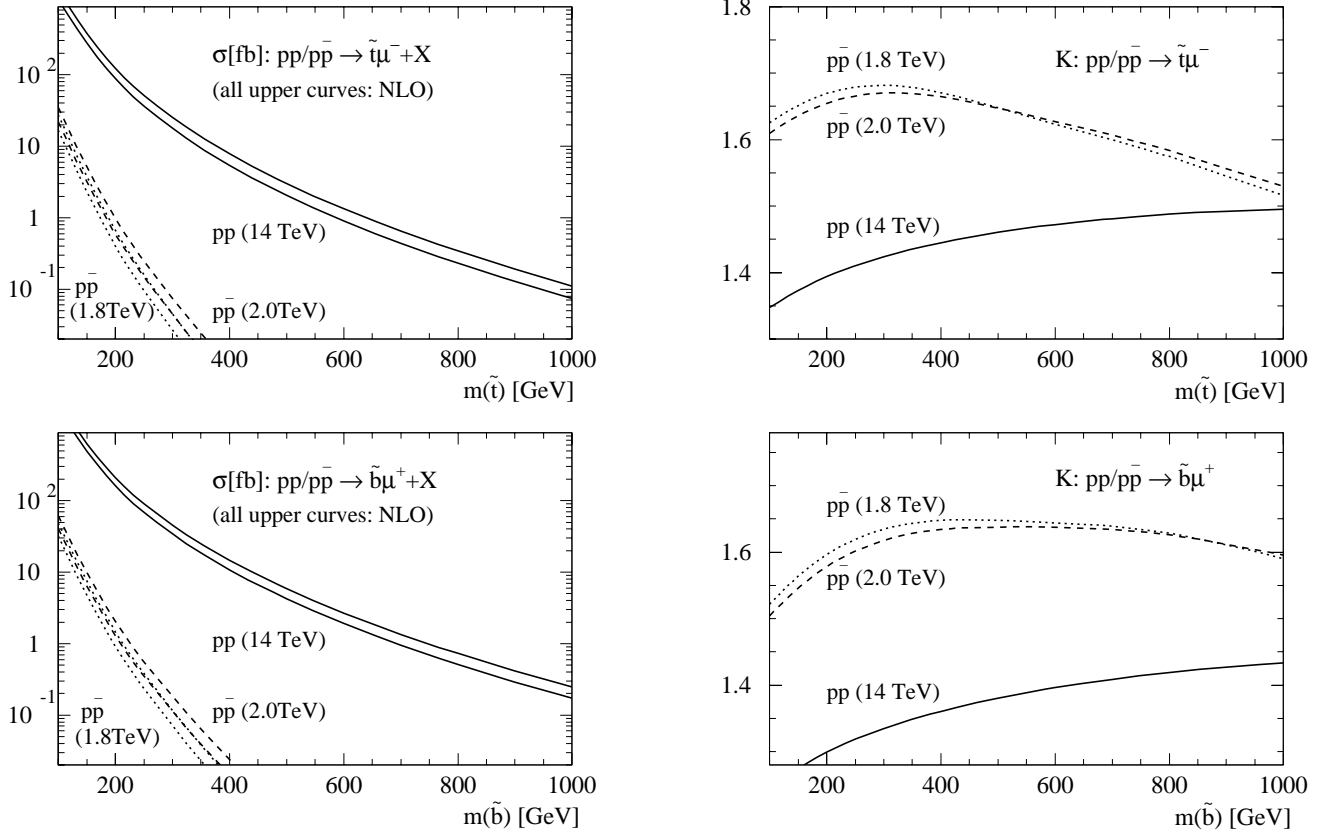
Figure 1. The impact of the next-to-leading order corrections for the associated production of a top (upper row) and bottom (lower row) squark as a function of the particle mass, shown for the leading order and next-to-leading order cross sections (left) and the K factor (right). All scales are set to their central value, *i.e.* the squark mass. The SUSY spectrum as a function of the stop (sbottom) mass is given in the text. The squark mixing angle, as well as the R parity violating coupling λ'_{SUSY}, are set to unity.

leptons. In order to avoid strong bounds on λ'_{SUSY} and have a clear and unsuppressed signal, we can for example identify $\lambda'_{\mathrm{SUSY}} = \lambda'_{231}$, which implies that the initial quark has down flavor and the associated lepton is a muon. In the case of sbottoms, the Lagrangean allows their production in association either with a charged lepton or with a neutrino. In that case we, for example, identify $\lambda'_{\mathrm{SUSY}} = \lambda'_{213}$, which corresponds to an incoming up quark and again a muon in the final state. We emphasize that all higher order results are the same for sbottom–lepton and sbottom–neutrino production, since the SUSY-QCD corrections do not see the charge of the lepton.

The complete set of $\mathcal{O}(\lambda'^2_{\mathrm{SUSY}}\alpha_s^2)$ corrections includes virtual gluon and gluino diagrams and real gluon emission diagrams, as well as the (crossed) processes $gg \to \tilde{t}_1 \ell \bar{q}$, $q\bar{q} \to \tilde{t}_1 \ell \bar{q}$, and $qq \to \tilde{t}_1 \ell q$. Several flavor combinations of the incoming and outgoing quarks in these crossed processes are possible. The treatment of heavy supersymmetric states is reviewed in Section III. We renormalize the couplings λ'_{SUSY} and α_s in the $\overline{\mathrm{MS}}$ scheme, while the final state stop (sbottom/leptoquark) mass and the squark mixing angle are renormalized using a generalized on–shell scheme, which includes a running mixing angle, evaluated with the stop mass as the fixed renormalization scale. Using this scheme for the mixing angle [10] is known to lead to problems for weak $\mathcal{O}(\alpha)$ corrections [11], but it is certainly best suited for QCD corrections.

The observable cross section in terms of a running mixing angle can, of course, be linked to the same observable in any other renormalization scheme, and the numerical effect has been shown to be negligible for a large set of renormalization schemes [10]. The purely gluon or quark induced subprocesses lead to a double counting with pair production in the case of an intermediate on–shell stop (sbottom/leptoquark): $gg, q\bar{q} \to \tilde{t}_1 \tilde{t}_1^* \to \tilde{t}_1(\ell\bar{q})$. The on–shell contributions of these processes are usually evaluated as stop pair production with a subsequent R-parity violating decay. The off–shell contributions, however, should be part of the next-to-leading order corrections to the associated stop and lepton production. We, therefore, split the contributions into these two parts, using the small width approximation, and explicitly subtract the on–shell contribution coming from the squark pair production in the corresponding phase space points [12].

We display in Fig. 1 the effect of the next-to-leading order SUSY-QCD corrections to the associated stop–lepton and sbottom–lepton productions. While the actual value of the cross sections depends on the numerical value of the R-parity violating coupling λ'_{SUSY} and on the squark mixing angle, both of them drop out for the K factor, which is defined consistently as $K = \sigma_{\mathrm{NLO}}/\sigma_{\mathrm{LO}}$. As expected, there is hardly any difference for the Tevatron Run I and Run II results. Because the limited energy of the Tevatron makes it increasingly harder to radiate additional jets, we see that the next-to-leading order correction becomes smaller for heavier squark masses. The K factor at the LHC looks qualitatively different. While the center-of-mass energy is large compared to the squark mass, a sizeable fraction of the next-to-leading order cross section comes from two initial state gluons. For small squark masses it is likely that the purely gluonic initial state first produces an on-shell squark pair. This contribution we subtract, which automatically yields a small K factor. For larger squark masses more of the intermediate stops will be off–shell and, thereby, contribute to the K factor. On a very moderate level the difference between the K factors for the Tevatron Run I and Run II already shows the same effect which we see very clearly at the LHC. We emphasize that the smaller K factor for the LHC, as seen in Fig. 1 is entirely a function of the squark mass and not a feature of the next-to-leading order corrections at the LHC.

In Fig. 1 we also observe a difference in the cross section and in the K factor between stop and sbottom production. For the purely QCD corrections this is caused by the exchange of an incoming down quark in the stop case by an incoming up quark in the sbottom case. The supersymmetric corrections involving virtual gluinos require virtual quarks in the loops which match the flavor of the final state squark. Since the top mass is essentially of the same order of the supersymmetric mass scale and the bottom mass is very much smaller, this effect becomes visible. We note, however, that the behavior of the K factor as a function of the squark mass at the Tevatron and at the LHC is the same for final state stops and sbottoms, as one would expect from the arguments given above.

For production cross sections involving strongly interacting final states, the appropriate measure for the theoretical uncertainty coming from higher order corrections is the renormalization and factorization scale dependences. We show these scale dependences in Fig. 2. It is common to identify both of these scales, however, this can lead to systematic cancellations in the variation of the cross section with this common scale [13,5]. This, in turn, would yield a significant underestimate of the theoretical uncertainty. The scale dependence for the associated stop–lepton production is shown for a stop mass of 200 GeV and the corresponding SUSY spectrum. If one defines a theoretical error of the leading order cross section corresponding to the variation $Q = m_{\tilde{t}_1}/4 \cdots 4m_{\tilde{t}_1}$ of the renormalization scale dependence, the theoretical uncertainty at the Tevatron (LHC) is of the order of 20%. Notice that, in leading order, this uncertainty comes from the variation of $\alpha_s(Q_R)$. To next-to-leading order this error is reduced to 15% for the Tevatron, and slightly less for the LHC.

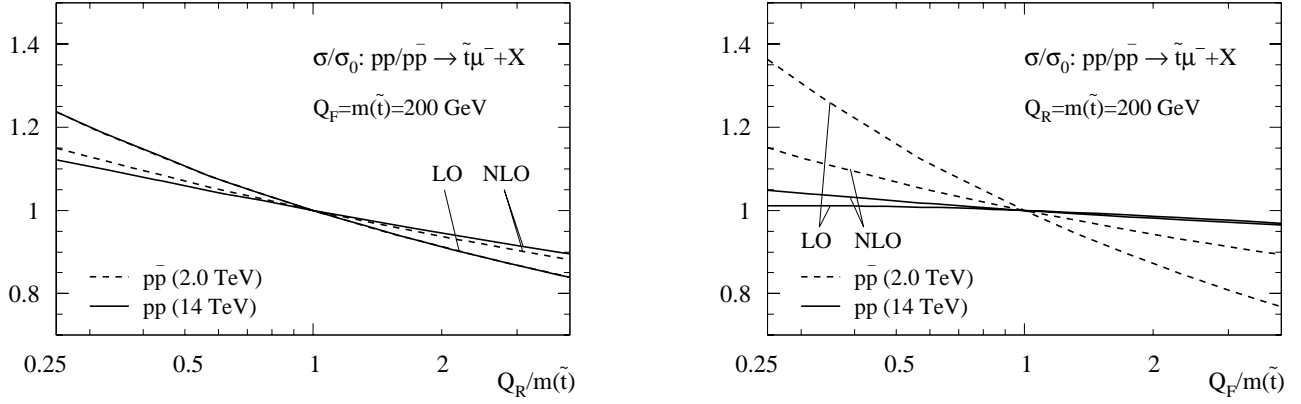In contrast, the factorization scale dependence of the leading order cross section is not identical for

Figure 2. Dependence of the total cross sections on the renormalization (left) and factorization (right) scales. The other respective scale is fixed to the central scale. The stop mass is set to 200 GeV and fixes the central value for the scales as well as the SUSY spectrum.

the Tevatron and the LHC. This is due to the fact that at the Tevatron the stops have to be produced right at threshold while at the LHC they can be produced through partons with considerably lower parton momentum fraction x. The leading order factorization scale dependence of this process at the LHC is extraordinarily mild and, in comparison with other similar processes [12,5], artificially flat. This is an effect of the choice of the renormalization scale and a cancellation involving a combination of logs $\log Q_R^2/m_{\tilde{t}_1}^2 \times \log Q_F^2/m_{\tilde{t}_1}^2$ [13]. The leading order factorization scale dependence actually changes the sign of the slope from positive to negative values when the stop mass increases from 100 to 500 GeV. Evaluating the scale dependence of the total cross section at the LHC for a stop mass of 200 GeV is very close to the turnover point, at which the factorization scale is actually flat.

Both leading order scale uncertainties add up to some 65% for the Tevatron and 40% for the LHC. The next-to-leading order calculations reduce this uncertainty significantly to $\sim 25\%$ for both colliders. Since there is no cancellation in the cross section between the two scale variations at next-to-leading order, we can obtain the same estimates by varying only the identified scale $Q = Q_R = Q_F$. These percentages are, of course, only a rough theoretically motivated estimate. In particular, the leading order uncertainty is more dependent on the powers of α_s in the cross section than on the actual process [12] and, therefore, has to be taken with a grain of salt. A comparison with the actual K factor given in Fig. 1 also shows that the estimated leading order uncertainty for the LHC is indeed too small, while the leading order error band covers the next-to-leading order curve well for the Tevatron. Furthermore, to next-to-leading order different supersymmetric production processes give very consistent uncertainty estimates [5,6,12,13]. Since there is no physics reason why the Tevatron and the LHC cross section should behave any differently as far as the theoretical uncertainty is concerned, it is a good check to see that at the next-to-leading order level the scale dependence is indeed similar for both experiments. Moreover, the next-to-leading error bands agrees well with what one would expect from similar processes.

We present in Fig. 3 the effect of the next-to-leading order contributions on the lepton transverse momentum and rapidity distributions. We generally observe that the next-to-leading order corrections change the shape of the lepton transverse momentum and rapidity distributions only at a negligible level. Looking into the transverse momentum distributions in more detail, we see that, while for the Tevatron the lepton
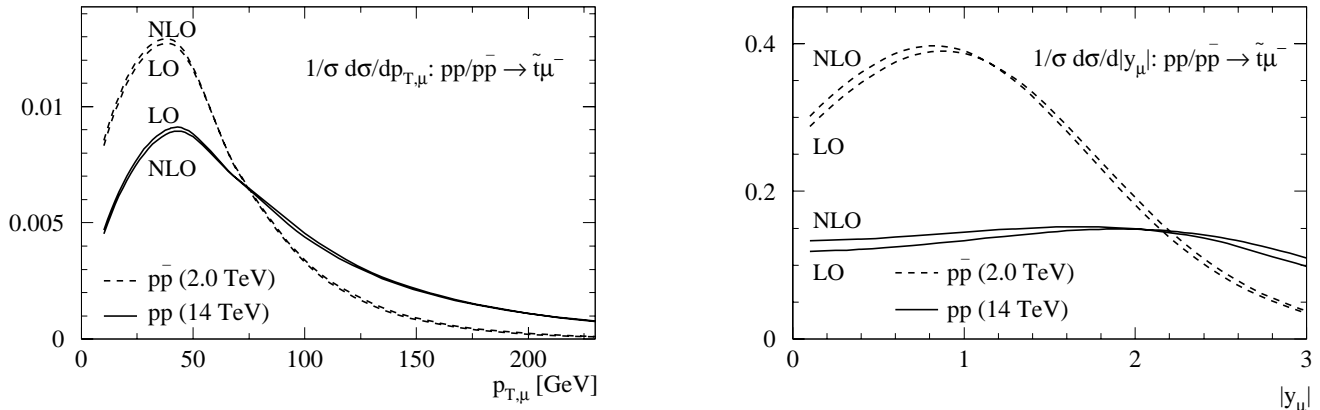
Figure 3. The normalized charged lepton transverse momentum (left) and rapidity (right) distributions at the Tevatron and at the LHC. The stop mass is set to 200 GeV and fixes the SUSY spectrum as well as the scale.

is expected to be slightly softer after including the next-to-leading order corrections, at the LHC the lepton tends to be harder. The reason is again the ratio of the stop mass and the collider energy: at the Tevatron all stops are produced very close to threshold with little transverse momentum. Therefore, there is very little energy available to produce additional hard jets; soft and collinear radiation will dominate the next-to-leading order effects and reduce the available partonic energy for the hard process. This renders the event softer altogether. In contrast, at the LHC there is more energy available to actually produce hard jets in initial and final state radiation. While initial state radiation still lowers the energy available for the hard process, final state radiation off the stop has to be balanced by the lepton. These two effects cancel to a very large degree but yield a slightly harder lepton spectrum. Because the leading order and next-to-leading order distributions are very similar, it is a sufficient approximation to rescale the leading order distributions by the constant K factor obtained from the total cross section.

The salient feature of the right panel of Fig. 3 is that the lepton rapidity distribution does not have its maximum at zero rapidity. In other words, the lepton prefers to be boosted. Looking at the initial state, we see that for the Born process the two incoming partons are one gluon and one valence quark. On average the latter will have the larger parton energy. This will boost the event altogether into the direction of the valence quark, which in turn, boosts the lepton into the forward direction. As can be seen in Fig. 3, the lepton is even farther forward at the LHC than it is at the Tevatron, which happens because at the LHC a larger fraction of events actually probes smaller parton momentum fractions x of the gluon and leads to larger momentum imbalance of the initial state. The next-to-leading order prediction softens the bias toward forward leptons at both colliders because of the final state radiation yields another hard central jet radiated from the stop, which has to be balanced by the lepton.

III. DECOUPLING OF HEAVY SUPERSYMMETRIC PARTICLES

While most of the calculation described in this paper has been done in the framework of R-parity violating supersymmetry, it is well known that the same kind of signal can be generated simply extending the Standard Model by a single scalar leptoquark [14]. This scalar leptoquark would correspond in supersymmetry to the lightest squark, which can be either the lighter stop ($\tilde{t}_1$) or the lighter sbottom ($\tilde{b}_1$). The
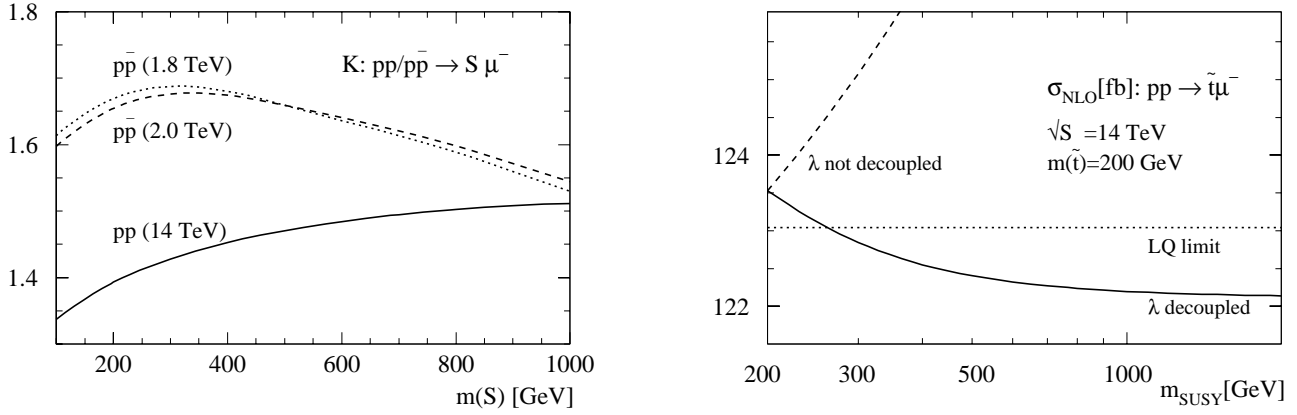
Figure 4. Left: the impact of the next-to-leading order QCD corrections on the leptoquark production cross sections. All input parameters are identical to the right panel of Fig. 1, the leptoquark coupling is derived from the stop case, *i.e.* the scalar leptoquark couples to an incoming down quark and an outgoing muon. Right: the decoupling behavior of the supersymmetric single stop production cross section at next-to-leading order. The dashed curve we compute without decoupling the heavy supersymmetric particles from the running of the R parity violating coupling $\lambda'_{\rm SUSY}$, as described in Eq. (6).

R-parity violating coupling $\lambda'_{\rm SUSY}$ corresponds to the scalar leptoquark coupling $\lambda'_{\rm LQ}$. In other words, if we decouple all additional strongly interacting supersymmetric particles, we should recover the Standard Model with an additional scalar leptoquark in leading order as well as in next-to-leading order QCD.

The decoupling of heavy strongly interacting states is slightly complicated by the contribution of these states to the running of $\overline{\rm MS}$ quantities and their identification with the corresponding Standard Model quantities. As an example we recall the treatment of heavy states in the evolution of the strong coupling $\partial \alpha_s / \partial \log \mu_R^2 = -\beta_\alpha \alpha_s / (4\pi)$. The current extraction of the strong coupling from data explicitly uses only five quark flavors, while the renormalization of the same quantity in the Standard Model involves all six quarks. As a matter of fact, the $\overline{\rm MS}$ α_s renormalization in the MSSM leads to the complete one-loop beta function, which does not exhibit the decoupling of the heavy states at low energy

$$\beta_\alpha = \beta_\alpha^{\rm SM} + \beta_\alpha^{\rm SUSY} = \left(\frac{11}{3} N - \frac{2}{3} n_f \right) + \left(-\frac{2}{3} N - \frac{1}{3} n_s \right) \tag{3}$$

The supersymmetric contribution comes from gluino and from squark loops; in the MSSM the number of quark and squark flavors is $n_f = n_s = 6$. To match the measured α_s, we have to explicitly decouple the heavy particles from the running of α_s, *i.e.* we have to include an additional term in the next-to-leading order corrections [12]:

$$\frac{\Delta \alpha_s}{\alpha_s} = \frac{\alpha_s}{4\pi} \left[\frac{1}{6} \log \frac{\mu_R^2}{m_t^2} + \frac{n_f - 1}{3} \log \frac{\mu_R^2}{m_{\tilde{q}}^2} + \frac{1}{6} \log \frac{\mu_R^2}{m_{\tilde{t}_1}^2} + \frac{1}{6} \log \frac{\mu_R^2}{m_{\tilde{t}_2}^2} + \frac{N}{6} \log \frac{\mu_R^2}{m_{\tilde{g}}^2} \right] \tag{4}$$

This contribution explicitly cancels the contribution of the heavy particles to the running of α_s and ensures that the one-loop renormalization of α_s is governed by $\beta_\alpha^{\rm SM}$ with $n_f = 5$.

As expected, we find a completely analogous behavior for the running of the R-parity violating coupling: $\partial \lambda'_{\rm SUSY} / \partial \log \mu_R^2 = -\beta_\lambda \alpha_s / (4\pi)$. At the one-loop level it turns out that the beta function of $\lambda'_{\rm SUSY}$ vanishes.

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This is not a generic feature of the MSSM but a pure accident. For example, the R-parity violating coupling λ'' has a finite one-loop beta function in the MSSM as well as in the limit of heavy gluinos [5]. The beta function of λ'_{LQ} is well-known [15]:

$$\beta_\lambda = \beta_\lambda^{\mathrm{LQ}} + \beta_\lambda^{\mathrm{SUSY}} = \left(\frac{3}{2} C_F\right) + \left(-\frac{3}{2} C_F\right) \tag{5}$$

Since neither λ'_{LQ} nor λ'_{SUSY} are measured parameters (yet) we can in principle choose any definition as long as the calculation is consistent. However, we cannot use the constant coupling λ'_{SUSY} and naively decouple the heavy supersymmetric particles to obtain the scalar leptoquark theory: the naive supersymmetric corrections would be divergent with $\log m_{\mathrm{heavy}}^2$. Again we have to include a one-loop decoupling term to cancel this logarithmic divergence of the next-to-leading order cross section. Throughout the numerical analysis we include this divergent next-to-leading order decoupling contribution. Nevertheless, we also find that in the $\overline{\mathrm{MS}}$ scheme we are left with a finite threshold correction difference between λ'_{SUSY} and λ'_{LQ}. A complete decoupling term would read:

$$\frac{\Delta\lambda'_{\mathrm{SUSY}}}{\lambda'_{\mathrm{SUSY}}} = \frac{1}{2}\frac{\alpha_s}{4\pi}\left[3\,C_F \log\frac{\mu_R^2}{m_{\tilde g}^2} + C_F\right] . \tag{6}$$

Notice that, in contrast to all divergent decoupling corrections, this last term does not vanish in the case of a light supersymmetric spectrum $\mu_R \sim m_{\tilde t_1} \sim m_{\tilde g}$. For this reason we do *not* include it in our numerical analyses. Initially, we can see in the left panel of Fig. 4 that the next-to-leading order corrections for a stop–type leptoquark coupling to an initial state down quark behave exactly like the ones for supersymmetric stops. In the right panel we exhibit the explicit decoupling of the heavy supersymmetric state, for which we assume the gluino mass, light-flavor squark mass, and the heavier stop mass as being degenerate and equal to M_{SUSY}. It is interesting to notice the extremely large effect that takes place if you do not decouple λ. Once all issues described above are taken care of, the supersymmetric cross section approaches a decoupling limit. However, the finite difference given in Eq. (6) remains between the supersymmetric decoupling limit and the leptoquark next-to-leading order model.

For the sake of completeness, we would like to point out that a consistent treatment of heavy flavors is also required in other sectors of the Standard Model. The top and bottom Yukawa couplings are running $\overline{\mathrm{MS}}$ parameters with contributions from heavy supersymmetric particles. These have to be removed from the running in the limit of a heavy supersymmetric spectrum exactly the way we described it for α_s and λ'_{SUSY} [13]. In contrast, the stop mass and the stop mixing angle in this paper are defined in the on–shell scheme, which does not involve any running and which is based on an independent physics condition. Therefore, they do not pose any problems for a decoupled supersymmetric spectrum.

IV. DISCUSSION AND OUTLOOK

The single production of stops (sbottoms/leptoquarks) has the potential of extending the reach of colliders to discover or rule out their existence [3–5]. As we have shown, the inclusion of higher order SUSY-QCD corrections renders the theoretical predictions for these processes more stable with respect to variations of the renormalization and factorization scales. Moreover, these corrections enhance the total cross section by a factor up to 70% at the Tevatron and 50% at the LHC. This leads to a larger sensitivity to stops (sbottoms/leptoquarks).

To illustrate the effect of QCD next-leading-order corrections on the searches for new physics, let us consider the single production of leptoquarks. Using the Born expressions for the cross sections, the attainable limits on leptoquarks coupling to up (down) quarks and charged leptons is 310 (260) GeV and 2.9 (2.4) TeV for the Tevatron RUN II and LHC, respectively [7], assuming that the leptoquark decays exclusively into a charged lepton and a quark. Taking into account the K factors displayed in Fig. 4 (left panel) these limits read 350 (280) GeV and 3.1 (2.6) TeV for the RUN II and LHC respectively.

It is interesting to notice that the single production of stops or sbottoms give us the opportunity to direct measure their Yukawa coupling to lepton–quark pairs, which otherwise would only be possible indirectly through their effects on the Drell-Yan process [16].

ACKNOWLEDGEMENTS

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