Gluino production in some supersymmetric models at the LHC

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Abstract

In this article we review the mechanisms in several supersymmetric models for producing gluinos at the LHC and its potential for discovering them. We focus on the MSSM and its left-right extensions. We study in detail the strong sector of both models. Moreover, we obtain the total cross section and differential distributions. We also make an analysis of their uncertainties, such as the gluino and squark masses, which are related to the soft SUSY breaking parameters.

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1 Introduction

Although the Standard Model (SM) [1], based on the gauge symmetry $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ describes the observed properties of charged leptons and quarks it is not the ultimate theory. However, the necessity to go beyond it, from the experimental point of view, comes at the moment only from neutrino data. If neutrinos are massive then new physics beyond the SM is needed.

Supersymmetry (SUSY) or symmetry between bosons (particles with integer spin) and fermions (particles with half-integer spin) has been introduced in theoretical papers nearly 30 years ago [2]. Since that time there appeared thousands of papers. The reason for this remarkable activity is the unique mathematical nature of supersymmetric theories, possible solution of various problems of the SM within its supersymmetric extensions as well as the opening perspective of unification of all interactions in the framework of a single theory [3, 4, 5].
However supersymmetry seemed, in the early days, clearly inappropriate for a description of our physical world\(^1\), for obvious and less obvious reasons, which often tend to be somewhat forgotten, now that we got so accustomed to deal with Supersymmetric extensions of the Standard Model. We recall the obstacles which seemed, long ago, to prevent supersymmetry from possibly being a fundamental symmetry of Nature\(^2\) [6].

We know that bosons and fermions should have equal masses in a supersymmetric theory. However, it even seemed initially that *supersymmetry could not be spontaneously broken at all* which would imply that bosons and fermions be systematically degenerated in mass, unless of course supersymmetry-breaking terms are explicitly introduced “by hand”. As a result, this lead to the question:

- Is spontaneous supersymmetry breaking possible at all?

Until today, the spontaneous supersymmetry breaking remains, in general, rather difficult to obtain. Of course just accepting the possibility of explicit supersymmetry breaking without worrying too much about the origin of supersymmetry breaking terms, as is frequently done now, makes things much easier – but also at the price of introducing a large number of arbitrary parameters, coefficients of these supersymmetry-breaking terms. These terms essentially serve as a parametrization of our ignorance about the true mechanism of supersymmetry breaking chosen by Nature to make superpartners heavy. In any case such terms must have their origin in a spontaneous supersymmetry breaking mechanism.

However, much before getting to the Supersymmetric Standard Model, and irrespective of the question of supersymmetry breaking, the crucial question, if supersymmetry is to be relevant in particle physics, is:

- Which bosons and fermions could be related?

But there seems to be no answer since known bosons and fermions do not appear to have much in common – except, maybe, for the photon and the neutrino.

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\(^1\) The most physicists considered supersymmetry as irrelevant for “real physics”.

\(^2\) We are grateful to P. Fayet that called our attention, and sent us all these interesting information about the “early” day of the Supersymmetric extensions of the Standard Model, as well as sent us all the original articles.
May be supersymmetry could act at the level of composite objects, e.g. as relating baryons with mesons? Or should it act at a fundamental level, i.e. at the level of quarks and gluons? (But quarks are color triplets, and electrically charged, while gluons transform as an $SU(3)$ color octet, and are electrically neutral!)

In a more general way the number of (known) degrees of freedom is significantly larger for the fermions (now 90, for three families of quarks and leptons) than for the bosons (27 for the gluons, the photon and the $W^\pm$ and $Z$ gauge bosons, ignoring for the moment the spin-2 graviton, and the still-undiscovered Higgs boson). And these fermions and bosons have very different gauge symmetry properties! This leads to the question:

- How could one define (conserved) baryon and lepton numbers, in a supersymmetric theory?

Of course nowadays we are so used to deal with spin-0 squarks and sleptons, carrying baryon and lepton numbers almost by definition, that we can hardly imagine this could once have appeared as a problem.

Supersymmetry today is the main candidate for a unified theory beyond the SM. Search for various manifestations of supersymmetry in Nature is one of the main tasks of numerous experiments at colliders. Unfortunately, the result is negative so far. There are no direct indications on existence of supersymmetry in particle physics, however there are a number of theoretical and phenomenological issues that the SM fails to address adequately [7]:

- Unification with gravity; The point is that SUSY algebra being a generalization of Poincaré algebra [5, 8, 9]

$$\{Q_{\alpha}, \bar{Q}_{\dot{\alpha}}\} = 2\sigma_{\alpha,\dot{\alpha}}^m P_m. \quad (1)$$

Therefore, an anticommutator of two SUSY transformations is a local coordinate translation. And a theory which is invariant under the general coordinate transformation is General Relativity. Thus, making SUSY local, one obtains General Relativity, or a theory of gravity, or supergravity [10].

- Unification of Gauge Couplings; According to hypothesis of Grand Unification Theory (GUT) all gauge couplings change with energy. All
known interactions are the branches of a single interaction associated with a simple gauge group which includes the group of the SM. To reach this goal one has to examine how the coupling change with energy. Considering the evolution of the inverse couplings, one can see that in the SM unification of the gauge couplings is impossible. In the supersymmetric case the slopes of Renormalization Group Equation curves are changed and the results show that in supersymmetric model one can achieve perfect unification [11].

- Hierarchy problem; The supersymmetry automatically cancels all quadratic corrections in all orders of perturbation theory due to the contributions of superpartners of the ordinary particles. The contributions of the boson loops are cancelled by those of fermions due to additional factor \((-1)\) coming from Fermi statistic. This cancellation is true up to the SUSY breaking scale, \(M_{\text{SUSY}}\), since

\[
\sum_{\text{bosons}} m^2 - \sum_{\text{fermions}} m^2 = M_{\text{SUSY}}^2, \tag{2}
\]

which should not be very large (\(\leq 1 \text{ TeV}\)) to make the fine-tuning natural. Therefore, it provides a solution to the hierarchy problem by protecting the electroweak scale from large radiative corrections [12]. However, the origin of the hierarchy is the other part of the problem. We show below how SUSY can explain this part as well.

- Electroweak symmetry breaking (EWSB); The “running” of the Higgs masses leads to the phenomenon known as radiative electroweak symmetry breaking. Indeed, the mass parameters from the Higgs potential \(m_1^2\) and \(m_2^2\) (or one of them) decrease while running from the GUT scale to the scale \(M_Z\) may even change the sign. As a result for some value of the momentum \(Q^2\) the potential may acquire a nontrivial minimum. This triggers spontaneous breaking of \(SU(2)\) symmetry. The vacuum expectations of the Higgs fields acquire nonzero values and provide masses to fermions and gauge bosons, and additional masses to their superpartners [13]. Thus the breaking of the electroweak symmetry is not introduced by brute force as in the SM, but appears naturally from the radiative corrections.

SUSY has also made several correct predictions [7]:
• Supersymmetry predicted in the early 1980s that the top quark would be heavy [14], because this was a necessary condition for the validity of the electroweak symmetry breaking explanation.

• Supersymmetric grand unified theories with a high fundamental scale accurately predicted the present experimental value of $\sin^2 \theta_W$ before it was measured [15].

• Supersymmetry requires a light Higgs boson to exist [16], consistent with current precision measurements, which suggest $M_h < 200$ GeV [17].

Together these successes provide powerful indirect evidence that low energy SUSY is indeed part of correct description of nature.

Certainly the most popular extension of the SM is its supersymmetric counterpart called Minimal Supersymmetric Standard Model (MSSM) [3, 18, 19].

The first attempt to construct a phenomenological model was done in [20], where the author tried to relate known particles together (in particular, the photon with a “neutrino”, and the $W^\pm$’s with charged “leptons”, also related with charged Higgs bosons $H^\pm$), in a $SU(2) \otimes U(1)$ electroweak theory involving two doublet Higgs superfields now known as $H_1$ and $H_2^3$. The limitations of this approach quickly led to reinterpret the fermions of this model (which all have 1 unit of a conserved additive R quantum number carried by the supersymmetry generator) as belonging to a new class of particles. The “neutrino” ought to be considered as a really new particle, a “photonic neutrino”, a name transformed in 1977 into photino; the fermionic partners of the colored gluons (quite distinct from the quarks) then becoming the gluinos, and so on. More generally this led one to postulate the existence of new $R$-odd “superpartners” for all particles and consider them seriously, despite their rather non-conventional properties: e.g. new bosons carrying “fermion” number, now known as sleptons and squarks, or Majorana fermions transforming as an $SU(3)$ color octet, which are precisely the gluinos, etc.. In addition the electroweak breaking must be induced by a pair of electroweak Higgs doublets, not just a single one as in the SM, which requires the existence of charged Higgs bosons, and of several

Then called $S = H_1$ left-handed and $T = H_2^2$ right-handed [6].
neutral ones [3, 18]. We also want to stress that on reference [3] were introduced squarks and gluinos (color octet of Majorana fermions, which couple to squark/quark pairs within what is now known as Supersymmetric Quantum Chromodynamics (sQCD)), that is the main subject of this article.

The still-hypothetical superpartners may be distinguished by a new quantum number called $R$-parity, first defined in terms of the previous $R$ quantum number as $R_p = (-1)^R$, i.e. +1 for the ordinary particles and −1 for their superpartners. It is associated with a $Z_2$ remnant of the previous $R$-symmetry acting continuously on gauge, lepton, quark and Higgs superfields as in [3], which must be abandoned as a continuous symmetry to allow masses for the gravitino [18] and gluinos [21]. The conservation (or non-conservation) of $R$-parity is therefore closely related with the conservation (or non-conservation) of baryon and lepton numbers, $B$ and $L$, as illustrated by the well-known formula reexpressing $R$-parity in terms of baryon and lepton numbers, as $(-1)^{2S}(-1)^{3B+L}$ [22]. This may also be written as $(-1)^{2S}(-1)^{3(B-L)}$, showing that this discrete symmetry may still be conserved even if baryon and lepton numbers are separately violated, as long as their difference $(B - L)$ remains conserved, at least modulo 2.

The finding of the basic building blocks of the Supersymmetric Standard Model, whether “minimal” or not, allowed for the experimental searches for “supersymmetric particles”, which started with the first searches for gluinos and photinos, selectrons and smuons, in the years 1978-1980, and have been going on continuously since. These searches often use the “missing energy” signature corresponding to energy-momentum carried away by unobserved neutralinos [3, 22, 23]. A conserved $R$-parity also ensures the stability of the “lightest supersymmetric particle”, a good candidate to constitute the non-baryonic Dark Matter that seems to be present in the Universe.

Massive neutrinos can also be naturally accommodated in $R$-parity violating supersymmetric theories, in which neutrinos can mix with neutralinos so that they acquire small masses [24, 25]. However, the phenomenological bounds on $B$ and/or $L$ violation [25, 26] can be satisfied by imposing $B$ as a symmetry and allowing the lepton number violating couplings to be large enough to generate Majorana neutrino masses.

However, the minimalistic extension of the MSSM is to introduce a gauge singlet superfield $\tilde{N}$, this model is called “Next Minimal Supersymmetric Standard Model” (NMSSM) [8, 27].

It is mainly motivated by its potential to eliminate the $\mu$ problem of the
MSSM [28], where the origin of the the $\mu$ parameter in the superpotential

$$W_{\text{MSSM}} = \mu H_1 H_2$$

is not understood. For phenomenological reasons it has to be of the order of the electroweak scale, while the “natural” mass scale would be of the order of the GUT or Planck scale. This problem is evaded in the NMSSM where the $\mu$ term in the superpotential is dynamically generated through the superpotential

$$W_{\text{NMSSM}} = \lambda \hat{H}_1 \hat{H}_2 \hat{N} - \frac{1}{3} k \hat{N}^3 + W_{\text{MSSM}}$$

The scalar component of $\hat{N}$ is the Higgs singlet with vacuum expectation value $x$. Therefore, this problem is evaded in the NMSSM where the $\mu$ term in the superpotential is dynamically generated through $\mu = \lambda x$ with a dimensionless coupling $\lambda$.

One of the simplest extensions of the SM that allows to naturally explain the smallness of the neutrino masses (without excessively tiny Yukawa couplings) consists in incorporating right-handed Majorana neutrinos, and imposing a see-saw mechanism \[4\] [30, 31] for the neutrino mass generation [28, 32], it is the “Minimal Supersymmetric Standard Model with three right handed neutrinos” (MSSM3RHN) [9].

The introduction of three families of the right-handed neutrinos $N^i$ (where $i$ is flavor indice) brings two new ingredients to the standard model; one is a new scale of the Majorana masses for the right-handed neutrinos, and the other a new matrix for Yukawa coupling constants of these new particles. Thus, we have two independent Yukawa matrices in the lepton sector as in the quark sector. Therefore, this model can accommodate a see-saw mechanism, and at the same time stabilise the hierarchy between the scale of new physics and the electroweak (EW) scale [33].

Other very popular ones are Left-Right symmetric theories (LRM) [34]. The main motivations for this model are that it gives an explanation for the parity violation of weak interactions, provides a mechanism (see-saw) for generating neutrino masses, and has $B - L$ as a gauge symmetry. The model has many predictions one can directly test at a TeV-scale linear collider [35].

The interesting and important features [36] of this kind of models are:

\[4\] For the see-saw mechanism the first paper is [29]
1. It incorporates Left-Right (LR) symmetry [34] (on top of the quark-lepton symmetry mentioned above), which leads naturally to the spontaneous breaking of parity and charge conjugation [34, 37].

2. It incorporates a see-saw mechanism for small neutrino masses [30, 31].

3. It predicts the existence of magnetic monopoles [38].

4. It leads to rare processes such as $K_L \rightarrow \mu \bar{\mu}$ through the lepto-quark gauge bosons (with however a negligible rate for $M_{PS} \geq 10 GeV$) [34].

5. In the case of single-step breaking, predicts the scale of quark-lepton (and Left-Right) unification [39].

6. It allows naturally for $\Delta B = 2$ process of $n - \bar{n}$ oscillations (with however a negligible rate unless there are light diquarks in the $TeV$ mass region) [40].

7. Last but not least, it allows for implementation of the leptogenesis scenario, as suggested by the see-saw mechanism [41].

On the technical side, the left-right symmetric model has a problem similar to that in the SM: the masses of the fundamental Higgs scalars diverge quadratically. As in the SM, the Supersymmetric Left-Right model (SUSYLR) can be used to stabilize the scalar masses and cure this hierarchy problem.

On the literature there are two different SUSYLR models. They differ in their $SU(2)_R$ breaking fields: one uses $SU(2)_R$ triplets [42] (SUSYLR T) and the other $SU(2)_R$ doublets [43] (SUSYLRD).

Another, maybe more important raison d’etre for SUSYLR is the fact that they lead naturally to R-parity conservation [44]. Namely, Left-Right models contain a $B-L$ gauge symmetry, which allows for this possibility [45]. All that is needed is that one uses a version of the theory that incorporates a see-saw mechanism [30, 31] at the renormalizable level.

As we said before, the supersymmetric particles have not yet been detected in the present machines such as HERA and Tevatron. By another hand there are many interesting supersymmetric models in the literature as we shown above. Therefore, if SUSY is detected in the Large Hadron Collider (LHC), one of the next steps will be to discriminate among the different
SM extensions, scenarios and also to find the mass spectrum of the different particles (which can be obtained theoretically in different scenarios).

To discriminate among the several possibilities, it is important to make predictions for different observables and confront these predictions with the forthcoming experimental data. One important process which could be measured at the LHC is the gluino production.

The \( R \)-symmetry transformations act chirally on gluinos, so that an unbroken \( R \)-invariance would require them to remain massless, even after a spontaneous breaking of the supersymmetry! In the early days it was very difficult to obtain large masses for gluinos, since: i) no direct gluino mass term was present in the Lagrangian density; and ii) no such term may be generated spontaneously, at the tree approximation, since gluino couplings involve colored spin-0 fields.

On this case, gluinos remain massless, and we would then expect the existence of relatively light “\( R \)-hadrons” \([22, 23]\) made of quarks, antiquarks and gluinos, which have not been observed. We know today that gluinos, if they do exist, should be rather heavy, requiring a significant breaking of the continuous \( R \)-invariance, in addition to the necessary breaking of the supersymmetry.

A third reason for abandoning the continuous \( R \)-symmetry could now be the non-observation at LEP of a charged \textit{wino} – also called \textit{chargino} – lighter than the \( W^{\pm} \), that would exist in the case of a continuous \( U(1) \) \( R \)-invariance \([3, 20]\). The just-discovered \( \tau^- \) particle could tentatively be considered, in 1976, as a possible light wino/chargino candidate, before getting clearly identified as a sequential heavy lepton.

The \textit{gluino masses} result directly from \textit{supergravity}, through \( m_{3/2} \), which leads one to abandon continuous \( R \) for \( R \)-parity as already observed in 1977 \([18]\). Another way does not use SUGRA but generates \( m_{\text{gluino}} \) radiatively, using messengers quarks sensitive to the source of SUSY-breaking \([46]\), on this case the gluino masses are generated by radiative corrections involving a new sector of quarks sensitive to the source of supersymmetry breaking, that would now be called “messenger quarks” \(^5\).

Today, the gluinos are expected to be one of the most massive sparticles which constitute the Minimal Supersymmetric Standard Model (MSSM), and

\(^5\)Remember that the see-saw mechanism \([29]\) did not attract attention at the time, however the article \([46]\) on gluino masses discusses a see-saw mechanism for gluinos.
therefore their production is only feasible at a very energetic machine such as the LHC. Being the fermion partners of the gluons, their role and interactions are directly related with the properties of the supersymmetric QCD (sQCD).

The aim of this paper is twofold. The first one is to study the strong sector of some supersymmetric models and to show explicit that the Feynmann rules for the gluino production are the same in all supersymmetric extensions of the Standard Model here considered. After that, as the second aim of this article, we show predictions for the gluino production on these models at the LHC, for various SPS benchmark points. The outline of the paper is the following. In sections 2 and 3, we obtain the relevant Feynman rules of the strong sector from the MSSM and SUSYLR models, respectively. Moreover, we see that the Feynman rules are indeed the same in these models. In section 4 we consider the different scenarios for the relevant SUSY parameters, which lead to different gluinos and squark masses. In section 5 we consider gluino production in the studied models and present the relevant expressions, which are used to obtain the numerical results in section 6. Conclusions are summarized in section 7.

2 Minimal Supersymmetric Standard Model (MSSM).

In the MSSM [3, 18, 19], the gauge group is $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$. The particle content of this model consists in associate to every known quark and lepton a new scalar superpartner to form a chiral supermultiplet. Similarly, we group a gauge fermion (gaugino) with each of the gauge bosons of the standard model to form a vector multiplet. In the scalar sector, we need to introduce two Higgs scalars and also their supersymmetric partners known as Higgsinos (Our notation is given at [47]). We also need to impose a new global $U(1)$ invariance usually called $R$-invariance, to get interactions that conserve both lepton and baryon number (invariance). On this section we will derive the Feynman rules of the strong sector of this model.

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6The particle content and the Lagrangian of this model.
2.1 Interaction from $\mathcal{L}_{\text{Gauge}}$

We can rewrite $\mathcal{L}_{\text{Gauge}}$, see [8, 9], in the following way

$$\mathcal{L}_{\text{Gauge}} = \mathcal{L}_{\text{cin}} + \mathcal{L}_{\text{gaugino}} + \mathcal{L}_{\text{gauge}}^D. \tag{5}$$

where

$$\begin{align*}
\mathcal{L}_{\text{cin}} &= \mathcal{L}_{\text{cin}}^{SU(3)} + \mathcal{L}_{\text{cin}}^{SU(2)} + \mathcal{L}_{\text{cin}}^{U(1)}, \\
\mathcal{L}_{\text{gaugino}} &= \mathcal{L}_{\text{gaugino}}^{SU(3)} + \mathcal{L}_{\text{gaugino}}^{SU(2)} + \mathcal{L}_{\text{gaugino}}^{U(1)}, \\
\mathcal{L}_{\text{gauge}}^D &= \mathcal{L}_{\text{D}}^{SU(3)} + \mathcal{L}_{\text{D}}^{SU(2)} + \mathcal{L}_{\text{D}}^{U(1)}, \tag{6}
\end{align*}$$

the first part is given by:

$$\mathcal{L}_{\text{cin}}^{SU(3)} = -\frac{1}{4} G_{mn}^a G^{amn}, \tag{7}$$

with

$$G_{mn}^a = \partial_m g_n^a - \partial_n g_m^a - g_s f^{abc} g_n^b g_m^c, \tag{8}$$

g_s is the strong coupling constant and $f^{abc}$ are totally antisymmetric structure constant of the $SU(3)$ group. The second term can be rewritten as

$$\mathcal{L}_{\text{gaugino}}^{SU(3)} = -\tau \lambda_C^a \bar{\sigma}^m D_m \lambda_C^a, \tag{9}$$

where

$$D_m \lambda_C^a = \partial_m \lambda_C^a - g_s f^{abc} \lambda_C^a g_m^c. \tag{10}$$

The last term is given by

$$\mathcal{L}_{\text{D}}^{SU(3)} = \frac{1}{2} D_C^a D_C^a. \tag{11}$$

2.1.1 Gluon Self Interaction

These interactions are derived from Eqs.(7,8) which are the same as the usual QCD. Therefore, on this case, we get the same Feynman rules as those from the QCD.
2.1.2 Gluino–Gluino–Gluon Interaction

This interaction is got from Eqs.(9, 10) combining both equations to obtain

\[ \mathcal{L}_{\text{gaugino}} = \mathcal{L}_{\text{cin}} + \mathcal{L}_{\tilde{g}\tilde{g}g}, \]  

where

\[ \mathcal{L}_{\text{cin}} = i(\partial_m \bar{\lambda}_a C) \bar{\sigma}^m \lambda_a C, \]
\[ \mathcal{L}_{\tilde{g}\tilde{g}g} = i g_s f^{abc} \bar{\lambda}_a C \bar{\sigma}^m \lambda_b C g_m, \]  

the first term gives the cinetic term to gluino, while the last one provides the gluino-gluino-gluon interaction.

Considering the four-component Majorana spinor for the gluino, given by

\[ \Psi(\tilde{g}^a) = \left( \begin{array}{c} -i \lambda^a_C(x) \\ i \lambda^a_C(x) \end{array} \right), \]  

we can rewrite \( \mathcal{L}_{\tilde{g}\tilde{g}g} \) in the following way

\[ \mathcal{L}_{\tilde{g}\tilde{g}g} = \frac{i}{2} g_s f^{bac} \bar{\Psi}(\tilde{g}^a) \gamma^m \Psi(\tilde{g}^b) g_m. \]  

Owing to the Majorana nature of the gluino one must multiply by 2 to obtain the Feynman rule (or add the graph with \( \tilde{g} \leftrightarrow \bar{\tilde{g}} \))!

The equation above induce the following Feynman rule, for the vertice gluino-gluino-gluon, given at Fig.(1) and it is the same result as presented

Figure 1: Feynman rule for the vertice Gluino-Gluino-Gluon at the MSSM.
2.2 Interaction from $\mathcal{L}_{\text{Quarks}}$

The interaction of the strong sector are obtained from the following Lagrangians

\begin{align*}
\mathcal{L}_{qqg} &= g_s \bar{Q}_m T^a Q g^{am} + g_s \bar{u}_r \sigma_m \bar{T}^a u^c g^{am} + g_s \bar{d}_r \sigma_m \bar{T}^a d^c g^{am} \\
\mathcal{L}_{q\bar{q}g} &= i g_s \left( \bar{Q}_m T^a \gamma^m Q - \bar{Q}_m T^a \gamma^m \tilde{Q} \right) g^{am} + i g_s \left( \bar{u}_r T^a \gamma^m u^c - \bar{u}_r T^a \gamma^m \tilde{u}^c \right) g^{am} \\
&\quad + i g_s \left( \bar{d}_r T^a \gamma^m \tilde{d}^c - \bar{d}_r T^a \gamma^m \tilde{d}^c \right) g^{am}, \\
\mathcal{L}_{\tilde{q}\tilde{q}g} &= -i g_s \left( \bar{\tilde{Q}}_m T^a \gamma^m \tilde{Q} - \bar{\tilde{Q}}_m T^a \gamma^m \tilde{Q} \right) g^{am} + i g_s \left( \bar{\tilde{u}}_r T^a \gamma^m \tilde{u}^c - \bar{\tilde{u}}_r T^a \gamma^m \tilde{u}^c \right) g^{am} \\
&\quad - i g_s \left( \bar{\tilde{d}}_r T^a \gamma^m \tilde{d}^c - \bar{\tilde{d}}_r T^a \gamma^m \tilde{d}^c \right) g^{am}, \\
\mathcal{L}_{\tilde{q}\bar{q}g} &= -g_s^2 \bar{\tilde{Q}}_m T^a \gamma^m Q g^{bm} - g_s^2 \bar{\tilde{Q}}_m T^a \gamma^m \tilde{Q} g^{bm} - g_s^2 \bar{\tilde{Q}}_m T^a \gamma^m \tilde{Q} g^{bm}.
\end{align*}

Where $T^a_{rs}$ are the color triplet generators, then one must use

\begin{equation}
T^a_{rs} = -T^a_{sr},
\end{equation}

for the color anti-triplet generators.

2.2.1 Quark–Quark–Gluon Interaction

This interaction comes from the first Lagrangian given at Eq.(16), and can be rewritten as

\begin{align*}
\mathcal{L}_{qqg} &= g_s \bar{u}_r \sigma^m T^a_{rs} u^c + \bar{d}_r \sigma^m T^a_{rs} d^c + \bar{u}_r \sigma^m T^a_{rs} u^c + \bar{d}_r \sigma^m T^a_{rs} d^c g^{am},
\end{align*}

$u$ and $d$ are color triplets while $u^c$ and $d^c$ are color anti-triplets. We must recall that $r$ and $s$ are color indices. Using Eq.(17) we can rewrite Eq.(18) as

\begin{align*}
\mathcal{L}_{qqg} &= g_s \bar{u}_r \sigma^m T^a_{rs} u^c + \bar{d}_r \sigma^m T^a_{rs} d^c + \bar{u}_r \sigma^m T^a_{rs} u^c + \bar{d}_r \sigma^m T^a_{rs} d^c g^{am}.
\end{align*}

Now, if we take into account the four-component Dirac spinor of the quarks ($q = u, d$), given by

\begin{align*}
\Psi(q) &= \begin{pmatrix} q_L(x) \\ q^c_L(x) \end{pmatrix},
\end{align*}

13
we obtain the Feynman rule

\[ \mathcal{L}_{qqg} = -g_s \sum_{q=u, d} \bar{\Psi}(q_r) \gamma^m T^a_{rs} \Psi(q_s) g_m^a, \]  

(21)

for the vertice \( qqg \) drawn at Fig.(2) this results again, as expected, agree

![Feynman Diagram](image)

Figure 2: Feynman rule for the vertice Quark-Quark-Gluon at the MSSM.

with the results given at [8, 9, 19, 48].

The gluino-gluon interaction is like the quark-gluon interaction, compare Fig.(1) with Fig.(2) \(^7\) and this fact gives the idea of “R-hadrons” presented at [22, 23], mentioned in our introduction.

### 2.2.2 Squark–Squark–Gluon Interaction

This interaction is obtained from the second line given at Eq.(16). This term can be rewritten in the following way

\[ \mathcal{L}_{\tilde{q} \tilde{q} g} = \, i g_s (\tilde{u}_r T^a_{rs} (\partial_m \tilde{u}_s) - (\partial_m \tilde{u}_r) T^a_{rs} \tilde{u}_s + \tilde{d}_r T^a_{rs} (\partial_m \tilde{d}_s) - (\partial_m \tilde{d}_r) T^a_{rs} \tilde{d}_s 
+ \bar{\tilde{u}} c_r T^a_{rs} (\partial_m \bar{\tilde{u}} c_s) - (\partial_m \bar{\tilde{u}} c_r) T^a_{rs} \bar{\tilde{u}} c_s + \bar{\tilde{d}} c_r T^a_{rs} (\partial_m \bar{\tilde{d}} c_s) - (\partial_m \bar{\tilde{d}} c_r) T^a_{rs} \bar{\tilde{d}} c_s) g^{am}. \]  

(22)

Using the following identity

\[ \tilde{A} \partial_m B \equiv A (\partial_m B) - (\partial_m A) B \]  

(23)

\(^7\)This fact as first noticed at [3].
we can rewrite our Lagrangian in the following simple way
\[ L_{\bar{q}\tilde{q}g} = i g_s \sum_{q=u,d} (\bar{q}_L T_a^a \tilde{q}_L m \bar{q}_L - \bar{q}_R T_a^a \tilde{q}_R m \bar{q}_R) g^{am}. \] (24)

Note the relative minus sign between the terms with \( \bar{q}_L \) and \( \bar{q}_R \): This is due to the facts that \( \bar{q}_R \) are colour anti–triplets and the anti–colour generator given at Eq.(17). Here, we use a similar notation as given at [19], it means that \( \bar{q}^* \) creates the scalar quark \( \tilde{q} \), while \( \bar{q} \) destroys the scalar quark \( \tilde{q} \).

Including the generalization to six flavors, we can write
\[ L_{\bar{q}\tilde{q}g} = i g_s \sum_{q=u,d} \sum_{p=1}^{6} (\bar{q}_{Lpr} T_a^a \partial_m \tilde{q}_{Lps} - \bar{q}_{Rpr} T_a^a \partial_m \tilde{q}_{Rps}) g^{am} \] (25)

The corresponding Feynman rule we obtain from
\[ \bar{q}_i^* \partial^m \tilde{q}_i = i \left(k_i + k_j\right)^m \] (26)
where \( k_i \) and \( k_j \) are the four–momenta of \( \tilde{q}_i \) and \( \tilde{q}_j \) in direction of the charge flow. This relations give us the following Feynman rules given at Fig.(3) and we conclude that this results, as expected, agree with the references

Figure 3: Feynman rule for the vertice Squark-Squark-Gluon at the MSSM. [8, 9, 19, 48].

2.2.3 Squark-Squark-Gluon-Gluon Interaction

This interaction comes from the last line given at Eq.(16), which can be written as
\[ L_{\bar{q}\tilde{q}gg} = -g_s^2 (\bar{u}_r T_a^a T_b^b \tilde{u}_t + \bar{d}_r T_a^a T_b^b \tilde{d}_t + \bar{u}_r T_a^a \bar{T}_a^b \tilde{u}_c + \bar{d}_r T_a^a \bar{T}_a^b \tilde{d}_c) g_m g^{bm}. \]
By using the following formula valid for \( SU(3) \) generators

\[
T^a_{rs} T^b_{st} = \frac{1}{6} \delta_{ab} \delta_{rt} + \frac{1}{2} (d^{abc} + i f^{abc}) T^c_{rt},
\]

(28)

it allows us to rewrite our Lagrangian in the following way\(^{8}\)

\[
\mathcal{L}_{\tilde q\tilde q g g} = -\frac{g_s^2}{6} \sum_{q = u, d} \tilde q^*_r T^a_{rs} \tilde q^a_m g^{am} - g_s^2 (d^{abc} + i f^{abc}) \sum_{q = u, d} \tilde q^*_r T^c_{rt} \tilde q^a_m g^{bm},
\]

(29)

However \( f^{abc} g^{a}_m g^{bm} = 0 \) because \( f^{abc} \) is totally antisymmetric structure constant of the group \( SU(3) \) while \( g^{a}_m g^{bm} \) is symmetric ones. The Feynman rule is drawn at Fig.(4) and it is in agreement with [8, 9, 19, 48].

Figure 4: Feynman rule to Squark-Squark-Gluon-Gluon vertex at MSSM.

### 2.2.4 Gluino-Quark-Squark Interaction

This interaction is described by the third line at Eq.(16), writing this term as

\[
\mathcal{L}_{\tilde q\tilde q g g} = -\sqrt{2} g_s (\tilde u^a_r T^a_{rs} \tilde u^a_s \lambda^a_C - \tilde u^a_r T^a_{rs} \tilde u^a_s \lambda^a_C + \tilde d^a_r T^a_{rs} \tilde d^a_s \lambda^a_C - \tilde d^a_r T^a_{rs} \tilde d^a_s \lambda^a_C)
\]

(30)

\[^{8}\text{Including the generalization to six flavors, see Eq.(25).}\]
Using the Eqs.\((20,14)\) and the usual chiral projectors
\[
L = \frac{1}{2} (1 + \gamma_5), \quad R = \frac{1}{2} (1 - \gamma_5),
\]
we can rewrite our Lagrangian in the following way
\[
\mathcal{L}_{\tilde{q}qq} = -\sqrt{2}g_s \sum_{q=u,d} (\bar{\Psi}(\tilde{g}^a)L \Psi(q_r) T_{rs}^a \tilde{q}_{sL}^* + \bar{\Psi}(q_r) R T_{rs}^a \bar{\Psi}(\tilde{g}^a)\tilde{q}_{sL} - \bar{\Psi}(q_r) L T_{rs}^a \Psi(\tilde{g}^a) \tilde{q}_{sL} - \bar{\Psi}(\tilde{g}^a) R \Psi(q_r) T_{rs}^a \tilde{q}_{sR}).
\]
this equation give us the following Feynman rule, given at Fig.\((5)\), and again

\[
\begin{align*}
\text{Figure 5: Feynman rule for the vertice Gluino-Quark-Squarks at the MSSM.}
\end{align*}
\]

it is in concordance with \([8, 9, 19, 48]\).

Before considering the left-right models, we want to say that the color sector of the interesting models NMSSM and MSSM3RHN are the same as in the MSSM model, therefore the results presented above are still hold on these models.

\section{Supersymmetric Left-Right Model (SUSYLR)}

The supersymmetric extension of left-right models is based on the gauge group \(SU(3)_C \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}\). On the literature, as we said at introduction, there are two different SUSYLR models. They differ in their \(SU(2)_R\) breaking fields: one uses \(SU(2)_R\) triplets \([42]\) (SUSYLRRT) and the other \(SU(2)_R\) doublets \([43]\) (SUSYLRD). Some details of both models are
described at [47]. SUSYLR models have the additional appealing characteristics of having automatic R-parity conservation.

In this article, we are interested in studying only the strong sector. As this sector is the same in both models, SUSYLRT and SUSYLRD, the results we are presenting in this section hold in both models.

3.1 Interaction from $\mathcal{L}_{\text{Gauge}}$

We can rewrite $\mathcal{L}_{\text{Gauge}}$, as done in the MSSM case, in the following way

$$\mathcal{L}_{\text{Gauge}} = \mathcal{L}_{\text{cin}} + \mathcal{L}_{\text{gaugino}} + \mathcal{L}_{\text{gauge}}^D.$$  \hfill (33)

where

$$\mathcal{L}_{\text{cin}} = \mathcal{L}_{\text{cin}}^{SU(3)} + \mathcal{L}_{\text{cin}}^{SU(2)\, L} + \mathcal{L}_{\text{cin}}^{SU(2)\, R} + \mathcal{L}_{\text{cin}}^{U(1)},$$

$$\mathcal{L}_{\text{gaugino}} = \mathcal{L}_{\text{gaugino}}^{SU(3)} + \mathcal{L}_{\text{gaugino}}^{SU(2)\, L} + \mathcal{L}_{\text{gaugino}}^{SU(2)\, R} + \mathcal{L}_{\text{gaugino}}^{U(1)},$$

$$\mathcal{L}_{\text{gauge}}^D = \mathcal{L}_{\text{gauge}}^{SU(3)} + \mathcal{L}_{\text{gauge}}^{SU(2)\, L} + \mathcal{L}_{\text{gauge}}^{SU(2)\, R} + \mathcal{L}_{\text{gauge}}^{U(1)},$$ \hfill (34)

the first part is given by:

$$\mathcal{L}_{\text{cin}}^{SU(3)} = -\frac{1}{4} G^a_{mn} G^{amn},$$ \hfill (35)

with $G^a_{mn}$ is given by Eq.(8)

$$\mathcal{L}_{\text{gaugino}}^{SU(3)} = -i \lambda^a C \bar{\sigma}^m D_m \lambda^a,$$ \hfill (36)

while $D^C_n \lambda^a_C$ is defined at Eq.(10). The last term is given by

$$\mathcal{L}_{\text{gauge}}^D = \frac{1}{2} D^a_C D^a_C.$$ \hfill (37)

The Eqs.(35,36) are equal from the Lagrangian of the MSSM, given at Eqs.(7,9). Therefore the Feynman rules are the same as in the MSSM case. The Feynman rule of the interaction Gluino-Gluino-Gluon is given at Eq.(15).
3.2 Interaction from $\mathcal{L}_{\text{Quarks}}$

In terms of the doublets the Lagrangian of the strong sector can be rewritten as

$$\mathcal{L}_{qqg} = g_s \tilde{Q} \bar{\sigma}_m T^a Q g^{am} + g_s \tilde{Q} \bar{\sigma}_m \tilde{T}^a Q^c g^{am},$$

$$\mathcal{L}_{q\bar{q}g} = i g_s \left( \tilde{Q} T^a \partial_m \tilde{Q} - \tilde{Q} T^a \partial_m \tilde{Q} \right) g^{am} + i g_s \left( \tilde{Q} \tilde{T}^a \partial_m \tilde{Q} - \tilde{Q} \tilde{T}^a \partial_m \tilde{Q} \right) g^{am},$$

$$\mathcal{L}_{q\bar{q}g} = -i \sqrt{2} g_s \left( \tilde{Q} T^a \tilde{Q} \tilde{g}^a - \tilde{Q} \tilde{T}^a \tilde{Q} \tilde{g}^a \right) - i \sqrt{2} g_s \left( \tilde{Q} \tilde{T}^a \tilde{Q} \tilde{g}^a - \tilde{Q} \tilde{T}^a \tilde{Q} \tilde{g}^a \right),$$

$$\mathcal{L}_{q\bar{q}g} = g_s^2 \tilde{Q} T^a T^b \tilde{Q} g_m g^{bm} - g_s^2 \tilde{Q} T^a \tilde{T}^b \tilde{Q} g_m g^{bm}. \tag{38}$$

3.2.1 Quark–Quark–Gluon Interaction

This interaction is given by the first line given at Eq.(38). Using the doublets we can write our Lagrangian in the following way

$$\mathcal{L}_{qqg} = g_s (\bar{u}_r \sigma^m T^a_{rs} u_s + \bar{d}_r \sigma^m \bar{T}^a_{rs} d_s + \bar{u}_r \sigma^m \bar{T}^a_{rs} u^c_s + \bar{d}_r \sigma^m \bar{T}^a_{rs} d^c_s) g^{am}, \tag{39}$$

which is the same as of the MSSM, see Eq.(18), therefore the Feynman rule is again given by Eq.(21).

3.2.2 Squark–Squark–Gluon Interaction

This interaction comes from the second line at Eq.(38), and can be rewritten as

$$\mathcal{L}_{q\bar{q}g} = i g_s \left( \bar{u}_r T^a_{rs} \partial_m \bar{u}_s - (\partial_m \bar{u}_r) T^a_{rs} \bar{u}_s + \bar{d}_r T^a_{rs} \partial_m \bar{d}_s - (\partial_m d_r) T^a_{rs} d_s \right) \tag{40}$$

which is the same of the MSSM, see Eq.(22), and the Feynman rule is given by Eqs.(25,26), as expected.

3.2.3 Squark-Squark-Gluon-Gluon Interaction

This interaction is given by the last line of Eq.(38), and it is given by

$$\mathcal{L}_{q\bar{q}g} = -g_s^2 (\bar{u}_r T^a_{rs} T^b_{st} \bar{u}_t + \bar{d}_r T^a_{rs} T^b_{st} \bar{d}_t + \bar{u}_r \bar{T}^a_{rs} \bar{T}^b_{st} \bar{u}_t + \bar{d}_r \bar{T}^a_{rs} \bar{T}^b_{st} \bar{d}_t) g^{am} g^{bm}. \tag{41}$$

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which is the same as Eq.(27), and the Feynman rule is given at Eq.(29).

3.2.4 Gluino-Quark-Squark Interaction

This interaction is given by the third line of Eq.(38), and can be rewritten as

$$L_{\tilde{q}gg} = -\sqrt{2}g_s(\bar{u}_r T^a_{rs} \tilde{u}_s \lambda^a C - \bar{u}_r T^a_{rs} u_s \lambda^a C + \bar{d}_r T^a_{rs} \tilde{d}_s \lambda^a C - \bar{d}_r T^a_{rs} d_s \lambda^a C + \bar{c}_r T^a_{rs} \tilde{c}_s \lambda^a C - \bar{c}_r T^a_{rs} c_s \lambda^a C),$$

(42)

which is given by Eq.(30), and the Feynman rule is given by Eq.(32).

As a conclusion of sections 2 and 3, we have shown that the Feynman rules of the strong sector are the same in the following MSSM, NMSSM, MSSM3RHN and SUSYLR models.

4 Parameters

The “Snowmass Points and Slopes” (SPS) [49] are a set of benchmark points and parameter lines in the MSSM parameter space corresponding to different scenarios in the search for supersymmetry at present and future experiments (See [50] for a very nice review). The aim of this convention is reconstructing the fundamental supersymmetric theory, and its breaking mechanism, from the experimental data.

The different scenarios correspond to three different kinds of models. The points SPS 1-6 are Minimal Supergravity (mSUGRA) model, SPS 7-8 are gauge-mediated symmetry breaking (GMSB) model, and SPS 9 are anomaly-mediated symmetry breaking (mAMSB) model ([49, 50, 51]), see appendix A.

Each set of parameters leads to different masses of the gluinos and squarks, which are the only relevant parameters in our study, and we shown their values in Tab.(1). In this paper, all these ten possibilities will be considered in our predictions for gluino production.

5 Gluino Production

Gluino and squark production at hadron colliders occurs dominantly via strong interactions. Thus, their production rate may be expected to be con-
Table 1: The values of the masses of gluinos and squarks in the SPS scenarios.

| Scenario | $m_{\tilde{g}}$ (GeV) | $m_{\tilde{q}}$ (GeV) |
|----------|----------------------|----------------------|
| SPS1a    | 595.2                | 539.9                |
| SPS1b    | 916.1                | 836.2                |
| SPS2     | 784.4                | 1533.6               |
| SPS3     | 914.3                | 818.3                |
| SPS4     | 721.0                | 732.2                |
| SPS5     | 710.3                | 643.9                |
| SPS6     | 708.5                | 641.3                |
| SPS7     | 926.0                | 861.3                |
| SPS8     | 820.5                | 1081.6               |
| SPS9     | 1275.2               | 1219.2               |

In the present paper we study the gluino production in $pp$ collisions. We will study the following reactions

$$pp \rightarrow \tilde{g}\tilde{g}, \tilde{g}\tilde{q} + X,$$

(43)

where X is anything, in the proton–proton collisions at the LHC.

In order to make a consistent comparison and for sake of simplicity, we restrict ourselves to leading-order (LO) accuracy, where the partonic cross-sections for the production of squarks and gluinos in hadron collisions were calculated at the Born level already quite some time ago [52]. The corresponding NLO calculation has already been done for the MSSM case [53], and the impact of the higher order terms is mainly on the normalization of the cross-sections.
Figure 6: Feynman diagrams for gluino pair production: (a) quark-antiquark initial states, (b) gluon-gluon initial states.

Figure 7: Feynman diagrams for squark–gluino production.

the cross section, which could be taken into account here by introducing a K factor in the results here obtained [53].

The LO QCD subprocesses for single gluino production are gluon-gluon and quark-antiquark annihilation (\(gg \rightarrow \tilde{g}\tilde{g}\) and \(q\bar{q} \rightarrow \tilde{g}\tilde{g}\)) (shown in Fig. 6), and the Compton process \(qq \rightarrow \tilde{g}\tilde{q}\) (shown in Fig. 7). For double gluino production only the annihilation processes contribute. These two kinds of events could be separated, in principle, by analysing the different decay channels for gluinos and squarks [8, 9].

Incoming quarks (including incoming \(b\) quarks) are assumed to be massless, such that we have \(n_f = 5\) light flavours. We only consider final state
squarks corresponding to the light quark flavours. All squark masses are taken equal to $m_{\tilde{q}}$. \footnote{L-squarks and $R$-squarks are therefore mass-degenerate and experimentally indistinguishable.} We do not consider in detail top squark production where these assumptions do not hold and which require a more dedicated treatment \cite{54}.

The invariant cross section for single gluino production can be written as \cite{52}

$$E \frac{d\sigma}{d^3p} = \sum_{ij} \int_{x_{min}}^{1} dx_a f_i^{(a)}(x_a, \mu) f_j^{(b)}(x_b, \mu) \frac{x_a x_b}{x_a - x_\perp \left( \frac{\zeta + \cos \theta}{2 \sin \theta} \right)} d\hat{t} (ij \to \tilde{g}d), \quad (44)$$

where $f_{i,j}$ are the parton distributions of the incoming protons and $\frac{d\sigma}{d\hat{t}}$ is the LO partonic cross section \cite{52} for the subprocesses involved. The identified gluino is produced at center-of-mass angle $\theta$ and transverse momentum $p_T$, and $x_\perp = \frac{2p_T}{\sqrt{s}}$. The Mandelstam variables of the partonic reactions $ij \to \tilde{g}\tilde{g}, \tilde{g}q$ are then

$$\hat{s} = x_a x_b s,$$
$$\hat{t} = m_{\tilde{g}}^2 - x_a x_\perp s \left( \frac{\zeta - \cos \theta}{2 \sin \theta} \right),$$
$$\hat{u} = m_{\tilde{g}}^2 - x_b x_\perp s \left( \frac{\zeta + \cos \theta}{2 \sin \theta} \right). \quad (45)$$

Here

$$x_b = \frac{2v + x_a x_\perp s \left( \frac{\zeta - \cos \theta}{\sin \theta} \right)}{2x_a s - x_\perp s \left( \frac{\zeta + \cos \theta}{\sin \theta} \right)},$$
$$x_{min} = \frac{2v + x_\perp s \left( \frac{\zeta + \cos \theta}{\sin \theta} \right)}{2s - x_\perp s \left( \frac{\zeta - \cos \theta}{\sin \theta} \right)},$$
$$\zeta = \left( 1 + \frac{4m_{\tilde{g}}^2 \sin^2 \theta}{x_\perp s} \right)^{1/2},$$
$$v = m_d^2 - m_{\tilde{g}}^2, \quad (46)$$

where $m_{\tilde{g}}$ and $m_d$ are the masses of the final-state partons produced. The center-of-mass angle $\theta$ and the differential cross section above can be easily
written in terms of the pseudorapidity variable $\eta = -\ln \tan(\theta/2)$, which is one of the experimental observables\(^{10}\). The total cross section for the gluino production can be obtained from above upon integration. The corresponding partonic total cross sections for the subprocesses considered are well known and can be found at [52, 53].

6 Numerical Results

Here in this section we present our numerical results and plots about the gluino production at the LHC. Since the $pp$ CM energy $\sqrt{s} = 14$ TeV is several times larger than the expected gluino and squark masses, these particles might be produced and detected at the LHC.

In Fig.8 we present the LO QCD total cross section for gluino production at the LHC as a function of the gluino masses. We use the CTEQ6L [55], parton densities, with two assumptions on the squark masses and choices of the hard scale (curves). The sensitivity with the hard scale is also presented in the case $m_{\tilde{q}} = m_{\tilde{g}}$. We also want to stress that the behaviour of our curves are similar to ones presented at Chapter 12 on reference [9], where they use the CTEQ5L parton distribution on their calculation.

The search for gluinos and squarks (as well as other searches for SUSY particles) and the possibility of detecting them will depend on their real masses. We also show (points) in Fig.8 the numerical results for the LO gluino total cross section in all SPS scenarios, fixing the gluino and squark masses, taken from Tab.(1).

The results show a strong dependence on the masses of gluinos and squarks. In the model curves, we get a larger cross section in the degenerated mass case, which agrees with [9]. Most of the SPS points are close to the first curve, which can be easily understood by looking at Table 1.

To discriminate among the different scenarios, it is relevant to look into more detailed observables such as differential distributions. In Fig.9 we present the transverse momentum distributions for single gluino production at LHC energies. The results show a huge difference in the magnitude for different scenarios - SPS1a (mSUGRA) gives the bigger values, SPS9 (AMSB) the smallest one. The predictions for the points SPS5 and SPS6

\(^{10}\)Remember that since gluinos are heavy, their rapidity and pseudorapidity are not equal.
Figure 8: The total LO cross section for gluino production at the LHC as a function of the gluino masses. Parton densities: CTEQ6L, with two assumptions on the squark masses and choices of the hard scale (curves). The sensitivity with the hard scale is also presented in the case $m_{\tilde{q}} = m_{\tilde{g}}$. The points are the numerical results for the SPS points as explained in the text.

(both mSUGRA) are indistinguishable, since the gluinos and squarks masses are almost the same in these two mSUGRA scenarios, so gluino production is not a good process to discriminate between them. Regarding the magnitude of the cross section of gluino production, this process could be useful to discriminate among basically four kinds of SPS scenarios, namely:

(i) SPS 1a;
(ii) SPS 5, SPS 6, and SPS 4 (in fact, this point is a bit lower than the other two);
(iii) SPS 1b, SPS 2, SPS 3, SPS 7, SPS 8;
(iv) SPS 9.

Looking into the details of the predictions, namely the behavior of the cross section, can give us further information to discriminate among the (iii)-
Figure 9: The LO $p_T$ distributions for single gluino production at the LHC ($|\eta| < 2.5$) for the different SPS points [49, 50]. We use CTEQ6L parton densities, and $\mu^2 = m_{\tilde{g}}^2 + p_T^2$ as a hard scale.

scenarios. The $p_T$ dependency in these scenarios is not the same for most of these points (except SPS 1b and SPS 3, which are almost equal). The SPS 2 (mSUGRA) model prediction has clearly a steeper falloff at high very high $p_T$, and the reason for that is the much higher squark mass in this scenario. At moderate $p_T$ of about 200 GeV, the SPS 7 (GMSB) curve has the lowest normalization, that is because the gluino mass is higher in this scenario. However, because the squark mass is much lower than in SPS 2, the falloff at higher $p_T$ is less steep and in this region the curve tend to other (iii)-points. So, the gluino and squark masses are more important in different regions and there is an interplay between them which produces different behaviors in
Figure 10: The LO $p_T$ distributions for double gluino production at the LHC ($|\eta| < 2.5$) for the different SPS points [49, 50]. We use CTEQ6L parton densities, and $\mu^2 = m_{\tilde{g}}^2 + p_T^2$ as a hard scale.

different scenarios.

The gluino mass is important in all subprocesses, but the squark mass only contributes to the $q\bar{q}$ annihilation and the Compton-like process $qg \to \tilde{g}q$, because of the t-channel squark exchange in both subprocesses and, of course, the squark production in the latter one. Comparing these two processes, the Compton process is dominant. Therefore, the different behaviors of the cross section are mainly due to the Compton-like contribution.

A complementary analysis can be done by considering double gluino production, which is easy to obtain from the calculation above, by picking only the annihilation processes (Fig. 6). The results for double gluino production are shown in Fig. 10. Similarly to the previous case, the results show huge
differences in the magnitude of the cross section for different scenarios - SPS1a gives the bigger values, SPS9 the smallest one. Also, we find very close values for SPS1b, SPS3 (mSUGRA) and SPS7 (GMSB), which makes it difficult to discriminate between these mSUGRA and GMSB models. The same occurs for SPS5 and SPS6 (both mSUGRA). However, differently from the single gluino case, the $p_T$ dependencies are similar in all scenarios. Another difference is the magnitude of curves for the points SPS 2 and SPS 8, which can be clearly separated from the other (iii)-points described above. From all this, we conclude that both processes, single and double gluino production, are complementary and useful to make different discriminations among the SPS scenarios.

7 Conclusions

In this paper we have studied the color sector in some extensions of the SM, namely MSSM and SUSYLR, and some others. We have derived the Feynman rules for the strong sector, and have showed explicitly that they are the same in all these models. This happens because the strong sector is the same in all SUSY extensions considered.

There are several scenarios for SUSY breaking, within the SPS convention, which imply in different values for the masses of the supersymmetric particles. To find the correct SUSY breaking mechanism one has to consider different observables which could be measured when SUSY particles start to be detected. In this article, we analyse gluino production at the LHC.

Gluinos are color octet fermions and play a major role to understanding sQCD. Because of their large mass as predicted in several scenarios, up to now the LHC is the only possible machine where they could be found.

Because the Feynman rules for the strong sector are the same in all SM extensions considered, the gluino production cross section are indeed equal in these models. Our results are in this sense model independent, or conversely, gluino production is not a good process to discriminate among those SUSY extensions of SM.

Besides, our results depend on the gluino and squark masses and no other SUSY parameters. Since the masses of gluinos come only from the soft terms, measuring their masses can test the soft SUSY breaking approximations. We have considered all the SPS scenarios and showed the corresponding
differences on the magnitude of the production cross sections. From this it is easy to distinguish mAMSB from the other scenarios. However, it is not so easy to distinguish mSUGRA from GMSB depending on the real values of masses of gluinos and squarks. Gluino production cannot distinguish the two scenarios SPS1b and SPS7, provided the gluino and squark masses are almost similar in these two cases (the same occurs for SPS 5 and SPS 6). For the other scenarios, such discrimination can be done, especially if we consider both single and double gluino production as complementary processes.

Gluino production is not a good process to discriminate among the Supersymmetric Models, but can be helpful in determining the correct SUSY breaking scenario and to understanding supersymmetric quantum chromodynamics.

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A Tables of SPS Convention

On this appendix we present the parameters of the SPS convention, which are given at Table 2.

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| SPS Point | mSUGRA: $m_0$ | $m_{1/2}$ | $A_0$ | $\tan \beta$ |
|-----------|---------------|-----------|--------|-------------|
| 1a        | 100           | 250       | -100   | 10          |
| 1b        | 200           | 400       | 0      | 30          |
| 2         | 1450          | 300       | 0      | 10          |
| 3         | 90            | 400       | 0      | 10          |
| 4         | 400           | 300       | 0      | 50          |
| 5         | 150           | 300       | -1000  | 5           |

| mSUGRA-like: | $m_0$ | $m_{1/2}$ | $A_0$ | $\tan \beta$ | $M_1$ | $M_2 = M_3$ |
|--------------|-------|-----------|-------|--------------|-------|-------------|
| 6            | 150   | 300       | 0     | 10           | 480   | 300         |

| GMSB: | $\Lambda/10^4$ | $M_{\text{mes}}/10^3$ | $N_{\text{mes}}$ | $\tan \beta$ |
|-------|-----------------|-------------------------|------------------|-------------|
| 7     | 40              | 80                      | 3                | 15          |
| 8     | 100             | 200                     | 1                | 15          |

| AMSB: | $m_0$ | $m_{\text{aux}}/10^3$ | $\tan \beta$ |
|-------|-------|------------------------|-------------|
| 9     | 450   | 60                     | 10          |

Table 2: The parameters for the Snowmass Points and Slopes (SPS). On this table all the scenarios consider $\text{sign}\mu = +$ taken from [50].

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