GLUINOS LIGHTER THAN SQUARKS AND DETECTION AT THE LHC

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We explain how physics in the kaon sector proves useful in setting constraints in models where all the supersymmetric particles, except the gluino, are in the range $\sim (3 - 10) \text{ TeV}$. We also discuss the signals of these models at colliders, in particular at the LHC.

Introduction

Here we discuss two important features of scenarios where there are three families of heavy squarks, while the mass of the gluino is considerably lighter than the rest of the other supersymmetric particles. The two features we discuss are: (i) bounds that can be obtained from flavour changing neutral currents (FCNC) in the kaon sector and (ii) the signatures of these models at colliders. In section I we disambiguate our scenario from others existing in the literature. In section II we present the restrictions from kaon physics, that can prove useful in setting some bounds of these models, together with typical conditions set by an underlying family symmetry. Then in section IV, we discuss the signatures at colliders, specifically at the LHC.

1 Models with lighter gluinos than squarks

Since there are many models in the literature with gluinos lighter than squarks, we disambiguate the scenario that we consider here with those in the literature and specify here our interest on it. Gauge Anomaly Mediation Symmetry Breaking (GAMSB) scenarios set soft terms, at tree level, equal to zero at a high-scale before the symmetry is broken. Then, loop corrections induce non zero masses and, in particular, the masses of the gauginos, $M_i$, $i = 1, 2, 3$, are different (e.g. [1]). On the other hand, G2-MSSM models, based on supergravity, achieve at $M_P$, $m_0 = m_{3/2} = A_0$, while also achieving a split among the gaugino masses (e.g. [2]). Both models, however, typically have the common feature that $M_2 > M_1 > M_3$, which sets a lower mass to the gluino, while leaving heavier scalars. How large is the split between the squarks and the gluino, depends on the details of the particular models. However, in fact this kind of spectra is not difficult to obtain from effective supergravity scenarios coming from string compactifications where the overall modulus, and not the dilaton, gives the main contribution to gaugino masses [3]. On the other hand, generic split supersymmetry models [4] set, as a working condition, a split between scalars and fermions, but one has some freedom in setting the precise split, as long as an agreement with all the observable physical quantities is satisfied. Models based on an underlying $U(2)$ family symmetry (e.g. [5]) achieve a split between the masses of two families of heavier squarks and one lighter, together with a split of the gluino mass, which is also light.
The scenario that we consider here sets the following ranges for the masses of the gluino and the squarks, respectively,

\[ 900 \text{ GeV} \leq m_{\tilde{g}} \leq 2000 \text{ GeV}, \quad 3000 \text{ GeV} \leq m_{\tilde{q}} \leq 10000 \text{ GeV}. \]  

Spectra of this type agrees with all the observable physical quantities if \( 4 \lesssim \tan \beta \lesssim 20 \). The exact details of achieving electroweak symmetry breaking conditions require a small fine tuning, as is expected as soon as one starts increasing the scale of some supersymmetric particles. However, as mentioned before, these models could arise from some supergravity scenarios, so if the conditions in Eq. 1 are a result from boundary conditions at a high-scale, the fine-tuning would be controlled by the underlying scenario.

2 Limits from kaon physics

It is well known that if physics, beyond the SM, in the kaon sector would contribute at tree level, then the effective Hamiltonian measuring the \( \Delta S = 2 \) transitions would require a new physics scale of about \( \Lambda = 10^4 \text{ TeV} \). In the general MSSM, the contribution to \( \Delta S = 2 \) comes at the loop level, so the scale decreases, nevertheless it could be used to put constraints of FCNC processes from scenarios where only some particles are heavier (\( O(10) \text{ TeVs} \)) and other are much lighter. The values of \( \Delta m_K \) and the flavour violating parameter \( \epsilon_K \) can be used to set limits on the overall scale of the mass of the scalars and of course in the off-diagonal masses. \( \Delta m_K \) is proportional to the effective Hamiltonian of the \( \Delta S = 2 \) transitions, \( \Delta m_K = 2 \text{Re} \left( K^0 \left| H_{\text{eff}}^{S=2} \right| K^0 \right) \), while \( \epsilon_K \) is proportional to the imaginary part, \( |\epsilon_K| = |\text{Im} \left( K^0 \left| H_{\text{eff}}^{S=2} \right| K^0 \right)|/(\sqrt{2} \Delta m_K) \). The values of \( \Delta m_K \) experimentally and in the SM are respectively

\[ \Delta m_K = (3.483 \pm 0.0059) \times 10^{-15}, \quad \Delta m_{K(\text{SD})}^{\text{NNLO}} = (3.1 \pm 1.2) \times 10^{-15} \text{ GeV}, \]

which we can see that they agree at the one sigma C.L. so the use in physics beyond the SM is to set limits in the contributions that the new processes can give. On the other hand, it is well known that long distance effects contributing to \( \Delta m_K \) cannot be precisely calculated, so the best we can do is to use the most precise values of the short distance effects, which we have denoted as \( \Delta m_{K(\text{SD})}^{\text{NNLO}} \) and whose value has been calculated in 6. Despite the uncertainties due to the unknown long distance effects, one can still use \( \Delta m_K \) to set limits on the masses of the squarks because for generic values of \( (\delta_{XY})_{ij}, X, Y = L, R \) (left, right), the contributions to \( \Delta m_K \) can easily overcome values of \( \times 10^{-15} \). So effectively what we can do is to use then \( \Delta m_{K(\text{SD})}^{\text{NNLO}} \) and its corresponding uncertainty as \( |\Delta m_K^S| < \left| \frac{\Delta m_{K(\text{exp})}}{\Delta m_{K(\text{exp})}} - \Delta m_{K(\text{SD})}^{\text{NNLO}} \right| + 2(\sigma_{\Delta m_{K(\text{SD})}^{\text{NNLO}}} + \sigma_{\Delta m_{K(\text{exp})}}) \), to set a limit on the supersymmetric contribution, \( \Delta m_K^S \). The values of \( \epsilon_K \) experimentally and in the SM are respectively

\[ |\epsilon_K| = (2.228 \pm 0.11) \times 10^{-3}, \quad |\epsilon_K^{\text{NNLO}}| = (1.81 \pm 0.28) \times 10^{-3}, \]

although these values also agree within the 2\( \sigma \) C.L., and there are many uncertainties in all the quantities involved in the calculation of \( |\epsilon_K| \) and Hoelbling’s talk), we can use them to constraint the contributions from supersymmetric processes, because just as in the case of \( \Delta m_K \), they can easily overcome the experimental value. This time, however, we could still hope that supersymmetric contributions would prove useful in bringing closer the SM and the experimental values. The effective Hamiltonian, \( H_{\text{eff}} = \sum_i \delta_{ij} \mathcal{O}_i \eta_i + \sum_i \delta_{ij} \mathcal{O}_i \eta_i \), depends, through the Wilson coefficients \( C_i \) only on the flavour violating parameters \( (\delta_{XY})_{12} \), which can

\[ a^{\text{For the exact expressions of all the quantities mentioned here, please check}} \]

The details of the QCD evolution is going to be presented in 5.
be defined as \( (\delta^d_{XY})_{ij} = (\hat{m}_{XY})_{ij} / \sqrt{\hat{m}_{XX}^{d2}(\hat{m}_{YY}^{d2})_{ij}} \), where \( \hat{m}_{XY}^{d2} \) represent the effective mass terms appearing in the effective 6 \( \times \) 6 squared mass matrix giving masses to the six \( d \) squarks, in the basis where Yukawa couplings are diagonal. For the parameters \( (\delta^d_{LR})_{12} \), we have

\[
(\delta^d_{LR})_{12} = \frac{-a^d_{12} v_d}{\sqrt{\hat{m}_{LL}^{d2}(\hat{m}_{RR}^{d2})_{22}}} \approx A^d m_6 \epsilon r \leq 10^{-2},
\]

where of course the term proportional to the \( \mu \) term disappears since it is proportional to the Yukawa coupling. The term proportional to the trilinear term \( a^d = Y^d A^d \) can still be off-diagonal because proportionality to Yukawa couplings is just a boundary condition at \( M_F \) and whenever the terms \( A_f \neq 0 \), then \( a^d \) runs different than \( Y_i^f \). The parameter \( \epsilon \) in Eq. 4 refers to a FS parameter and so depending on the value of \( r \), they can saturate the value of \( O(10^{-2}) \). To get this maximum value, \( O(1) \) parameters in going from the third expression after the initial \( = \) sign were taken into account. The parameters \( (\delta^d_{LL})_{12} \) and \( (\delta^d_{RR})_{12} \) cannot be expressed as \( (\delta^d_{LR})_{12} \) because they are more model dependent than it, however in FS both, imaginary and real parts, typically do not exceed \( O(10^{-1}) \).

Another important motivation for this work is to reassess the importance of QCD contributions. Naively, one would think that at few TeVs, LO corrections would be enough, however as it was pointed out in [9], in the context of two heavy families of squarks, NLO effects can still have an important impact. In this context, where the three families of squarks are of the same order of magnitude, the case is not different. Then, according to the motivations above, what do here, as a complement of [8], is to

1. Assess the importance of considering NLO effects in setting the limits of the masses of \( d \) type squarks and the flavour violating parameters \( (\delta^d_{XY})_{12} \) for the case of three heavy squarks families and a light gluino.

2. Check that the values of squarks masses, in the range considered, and the real parts of \( (\delta^d_{XY})_{12} \) satisfy the bounds set by \( \Delta m_K \).

3. Put constraints on the imaginary parts of \( (\delta^d_{XY})_{12} \).

In Fig. 1 we show a plot of the mass of the gluino \( m_{\tilde{g}} \) versus the common mass scale of the three heavy squarks families. In this plot we have included: (i) a solid (blue) continue line corresponding to the NLO evolution, (b) a medium dashed (red) line corresponding to the LO approximation, (c) a small-dashed (orange) light line corresponding to the LO evolution using VIA approximation for bag parameters and (d) a dot-dashed (light blue) line corresponding to the results without QCD corrections. For comparison, we have additionally plotted a medium-dashed dark (black) line representing the NLO evolution of the case where only two families where decoupled at the scale \( m_{\tilde{q}} \), one is light and we have followed a similar QCD evolution to the three heavy family case. Apart from showing the importance of NLO QCD corrections, the curves in Fig. 1 show that for the scenarios in the present discussion, the proper QCD evolution, sets a safe limit on the masses of the squarks of the type \( d \) considered.

In Fig. 2 we present curves for fixed values of \( m_{\tilde{q}} \), plotting \( m_{\tilde{g}} \) against the limits on \( \sqrt{\text{Re}((\delta^d_{XY})_{12}(\delta^d_{XY})_{12})} \). Note then that, according to Eq. 4 at the scale of squarks of the type \( d \) of \( \leq O(10) \) TeV, the use of \( \Delta m_K \) proves relevant since for some cases, the limit set by it, is in fact close to the allowed value for \( (\delta^d_{RR})_{12} \). The same happens when contributions from \( (\delta^d_{RR})_{12} \) are present. For other cases, where only \( (\delta^d_{LR})_{12} \) and/or \( (\delta^d_{RL})_{12} \) contribute, we can check that these models give non-dangerous contributions to \( \Delta m_K \). Finally, in Fig. 3 we present curves for fixed values of \( m_{\tilde{q}} \), plotting \( m_{\tilde{g}} \) against the limits on \( \sqrt{\text{Im}(((\delta^d_{XY})_{12})(\delta^d_{XY})_{12}))} \).
Figure 1: Limits on the mass of the squarks of the type d, $m_{\tilde{q}}$, as a function of the mass of the gluino, $m_{\tilde{g}}$, for a case where only $\sqrt{|\text{Re}((\delta^{d}_{dLL})_{12}^{1} (\delta^{d}_{dRR})_{12}^{1})|}$ contributes, with a value equal to 0.1. The solid (blue) continue line corresponds to the NLO evolution, while the medium dashed (red) line corresponds to the LO approximation. Note that the difference between NLO and LO is of the $O(15\%)$, which proves the importance of the QCD NLO determination for these cases of heavy squarks. For the rest of the lines, please refer to the text.

Figure 2: Limits on $\sqrt{|\text{Re} ((\delta^{d}_{X'Y'})_{12}^{1} (\delta^{d}_{X'Y'})_{12}^{1})|}$, using $\Delta m_K$. In all cases only the flavour violating parameters shown in the vertical axis have been set different to zero. The limits on $\sqrt{|\text{Re}[(\delta^{d}_{RR})_{12}^{1}]|}$ and $\sqrt{|\text{Re}[(\delta^{d}_{RL})_{12}^{1}]|}$ can be obtained by interchanging $R \leftrightarrow L$ in the upper two panels, respectively. The two lower panels show, respectively, the results for $\sqrt{|\text{Re}[\delta^{d}_{LL}^{2} \delta^{d}_{RR}^{2}]|_{12}^{1}}$ and $\sqrt{|\text{Re}[\delta^{d}_{LR}^{2} \delta^{d}_{RL}^{2}]|_{12}^{1}}$. The different curves correspond, from bottom to top to $m_{\tilde{q}} = 4, \ldots, 10$ TeV.
As expected, the limits on these quantities, set by $\epsilon_K$, prove to be indeed relevant, since the limits lie above the maximum values that in principle FS can achieve, so these limits are a useful constraint for these scenarios.

3 Signatures at the LHC

This scenario features the same kind of signatures at colliders, as the models cited in section 1, which are mainly the following:

1. $\Delta m_\chi = m_{\chi^\pm} - m_{\chi^0} = O(100)$ MeV, where $m_{\chi^\pm}$ and $m_{\chi^0}$ are, respectively, the masses of the lightest chargino and the lightest neutralino. The order of magnitude of $\Delta m_\chi$ leads to the existence of soft pions.

2. 4 SM fermion final states as a result of the gluino decays

$$pp \rightarrow \tilde{g}\tilde{g}, \quad \tilde{\chi}_1^0 \tilde{\chi}_1^\pm, \quad \tilde{\chi}_1^\pm \tilde{\chi}_1^-$$

Decays of the secondary charginos/neutralinos

$$\tilde{g} \rightarrow \tilde{\chi}_2^0 tt, \quad \tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 W^+, \quad \tilde{\chi}_1^0 \rightarrow \tilde{\chi}_1^- \tilde{\nu}_e, \quad \tilde{\chi}_1^0 qq', \quad \tilde{\chi}_1^0 \rightarrow \tilde{\chi}_1^- \bar{u} + h.c. \quad (5)$$

3. Effective supersymmetric cross section at $\lesssim 2$ TeV, produced mainly by gluinos, that could be easily tested at the LHC.

The first of the items above occurs whenever there is a large $|\mu|$ parameter. At zeroth order, in fact $m_{\tilde{\chi}_1^0}^{(0)} - m_{\tilde{\chi}_1^0}^{(0)} = M_2$, at the first loop level they receive corrections from mixing with the Higgsino, then, at next order, there will be a mixing with the $U(1)_Y$ gaugino, enhanced by $\tan \beta$ and dependent on the difference of $M_1$ and $M_2$: $m_{\tilde{\chi}_1^0}^{(2)} - m_{\tilde{\chi}_1^0}^{(2)} = m_{\tilde{\chi}_1^0}^{(0)} \tan^2 \beta \frac{\mu}{(M_1 - M_2)\mu^2} \sin^2 2\beta$. This condition produces the decay of the charginos into neutralinos and soft-pions, $\chi^\pm \rightarrow \chi^0 \pi^\pm$, then there would be displaced vertices with a track of few centimeters long. However, current strategies of long-lived charged particles that make it all the way through the detector seem to evade the detection of the kind of decays generated by this kind of neutralinos (e.g. $[10]$). However, if experimental techniques can overcome the difficulties in detection, this could provide indeed an interesting discerning criteria for these models. For systematic studies of this signature, although in different models than the present, see for example $[11]$. The second of the signals

Figure 3: Limits on $\sqrt{\text{Im}[\delta_{12}^X]\delta_{12}^K}$, using $\epsilon_K$. The limits on $\sqrt{\text{Im}[\delta_{12}^R]\delta_{12}^R}$ and $\sqrt{\text{Im}[\delta_{12}^L]\delta_{12}^L}$ can be obtained by interchanging $R \leftrightarrow L$, respectively. Limits for the case where $\sqrt{\text{Im}[\delta_{12}^R]\delta_{12}^R}$ and $\sqrt{\text{Im}[\delta_{12}^L]\delta_{12}^L}$, just as in the real case, are approximately one order of magnitude smaller than the corresponding case where only one of either $\delta_{12}^R$ or $\delta_{12}^L$ are present. The different curves correspond, from bottom to top, to $m_\tilde{q} = 4, \ldots, 10$ TeV.
above, although not unique for the scenario discussed here, are interesting because they clearly signal a split of the supersymmetric scalars. The difficulty in detection of course arises from the fact that it would be difficult at the LHC to disentangle the QCD background from a possible signal. Finally, for the third of the items above, LHC is already setting limits on what the effective supersymmetric cross section of related models should be. We find that a typical supersymmetric cross section, produced only by gluinos of $m_{\tilde{g}} = 650$ GeV, is $\sigma_{\text{SUSY}} \approx 0.2$ pb, at a luminosity $L \leq 800 \text{pb}^{-1}$, and an energy, at the center of mass, of $\sqrt{s}$ TeV. For gluinos of $m_{\tilde{g}} = 850$ GeV, $\sigma_{\text{SUSY}}$ decreases to $\approx 0.04$ pb, to finally decrease to $\sigma_{\text{SUSY}} \approx 0.001$ pb for $m_{\tilde{g}} = 1$ TeV (check [12] for similar results in the context of the G2-MSSM models). In fact, ATLAS and CDF have already put limits on related scenarios, where they set the mass of all the supersymmetric particles at 10 TeV, except of course the gluino mass. They have improved the previous exclusion of 2010 of gluinos with $m_{\tilde{g}} = 580$ GeV. In particular, ATLAS (see M. D. Joergensen talk and [13]) has determined that gluinos with a mass lower than $m_{\tilde{g}} = 810$ GeV are excluded. As explained before, the scenario considered here, sets the values of the squarks of the type d below, so the limits on the gluino masses from the experimental cross sections already determined, serve just as an illustration of the capacity of the LHC in setting bounds for the gluino mass. A detailed study at the LHC will require specific details of the kind of models presented here.

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References

1. Tony Gherghetta, Gian F. Giudice, and James D. Wells. Phenomenological consequences of supersymmetry with anomaly induced masses. *Nucl.Phys.*, B559:27–47, 1999.
2. Bobby Samir Acharya and Konstantin Bobkov. Kahler Independence of the G2-MSSM. JHEP, 09:001, 2010.
3. A. Brignole, Luis E. Ibáñez, and C. Muñoz. Soft supersymmetry-breaking terms from supergravity and superstring models. In *Kane, G.L. (ed.): Perspectives on supersymmetry II* 244-268, 1997.
4. Andrew G. Cohen, D.B. Kaplan, and A.E. Nelson. The More minimal supersymmetric standard model. Phys.Lett., B388:588–598, 1996.
5. Riccardo Barbieri, Gino Isidori, Joel Jones-Perez, Paolo Lodone, and David M. Straub. U(2) and Minimal Flavour Violation in Supersymmetry. *Eur.Phys.J.*, C71:1725, 2011.
6. Joachim Brod and Martin Gorbahn. Next-to-Next-to-Leading-Order Charm-Quark Contribution to the CP Violation Parameter $\epsilon_K$ and $\Delta M_K$. *Phys. Rev. Lett.*, 108:121801, 2012.
7. Kenji Kadota, Gordon Kane, Joern Kersten, and Liliana Velasco-Sevilla. Flavour Issues for String-Motivated Heavy Scalar Spectra with a low Gluino Mass: the $G_2$-MSSM case. *Eur. Phys. J. C* 72 (2012) 204.
8. J. Kersten and L. Velasco-Sevilla. Flavour constraints on scenarios with two or three heavy squark generations. [arXiv:1207.3016].
9. R. Contino and I. Scimemi. The supersymmetric flavor problem for heavy first-two generation scalars at next-to-leading order. *Eur. Phys. J. C* 10, 347 (1999).
10. Emma Torro. Prospects for detecting long-lived supersymmetric particles with ATLAS. J.Phys.Conf.Ser., 171:012088, 2009.
11. C.H. Chen, Manuel Drees, and J.F. Gunion. A Nonstandard string / SUSY scenario and its phenomenological implications. *Phys.Rev.*, D55:330–347, 1997.
12. Daniel Feldman, Gordon Kane, Ran Lu, and Brent D. Nelson. Dark Matter as a Guide Toward a Light Gluino at the LHC. Phys.Lett., B687:363–370, 2010.
13. ATLAS Collaboration. Search for charged long-lived heavy particles with the ATLAS Experiment at the LHC. [ATLAS-CONF-2012-022].