

COLOR-OCTET SCALARS AT THE LHC*

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Elements of the phenomenology of color-octet scalars (sgluons), as predicted in the hybrid $N = 1/N = 2$ supersymmetric model, are discussed in the light of forthcoming experiments at the CERN Large Hadron Collider.

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1. Motivation

After many years of preparations and constructions, the Large Hadron Collider at CERN will start its operation soon. One of the most important questions the LHC should provide an answer to is related to the nature of the electroweak gauge symmetry breaking. In the Standard Model (SM) it is achieved by adding a doublet of Higgs fields and arranging the parameters such that the symmetry is spontaneously broken. Although the SM is extremely successful in describing the experimental data, there are many arguments suggesting that it cannot be the ultimate theory, *e.g.* the SM Higgs sector is unnatural — cannot explain why the electroweak scale is so small with respect to the Planck scale; the SM cannot account for the matter–antimatter asymmetry of the Universe, nor for the dark matter *etc.* In fact, the above shortcomings point towards a new physics at the TeV scale.

Among many propositions for solving part of these problems in the physics area beyond the Standard Model, supersymmetry (SUSY) is generally considered most elegant and respected. Not invented or designed to, it can accommodate or explain some of the outstanding problems of the SM: stabilizing the gap between electroweak and Planck scale, unifying the gauge couplings, inducing radiative electroweak symmetry breaking. It

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provides a candidate for dark matter (DM), offers new ideas on the matter–antimatter asymmetry, while its unique mathematical structure can provide a link to physics at the GUT/Planck scale and, in local form, paving the path to gravity.

In the simplest $N = 1$ supersymmetric extension of the SM, each SM particle is paired with a sparticle that differs in spin by half a unit. Since none of the sparticles has been seen so far, the supersymmetric Lagrangian must be supplemented by SUSY breaking terms that keep unseen superpartners out of experimental reach while retaining renormalizability and maintaining perturbatively stable hierarchy of scales. Experimental constraints, mainly from flavor and Higgs physics, limit the allowed parameter space and play an increasingly restrictive role in building models of SUSY breaking.

However, the successes of supersymmetry do not rely on its simplest realization. In fact, non-minimal realizations may ameliorate the flavor problem. For example, Dirac gauginos (in contrast to Majorana in the MSSM) forbid some of the couplings and often lead to additional suppression of contributions from loops with gauginos in flavor-changing processes. Such scenarios can be based on D-term supersymmetry breaking models [1,2] or continuous R-symmetries [3].

A Dirac gaugino requires additional fermionic degrees of freedom. They can be provided by adding a chiral super-multiplet in the adjoint representation of the gauge group [4]. Here we will present elements of the phenomenology of the scalar partner of the Dirac gluino (sgluon), as worked out in a recent paper [5] (and confronted with Ref. [6]).

2. Short introduction to the $N = 1/N = 2$ hybrid model

In the MSSM gluinos are Majorana fields with two degrees of freedom to match gluons in the color-octet vector super-multiplet. To provide the two additional degrees of freedom for Dirac fields, the usual $N = 1$ gluon/gluino vector super-multiplet $\hat{g} = \{g_\mu, \tilde{g}\}$ may be paired with an additional $N = 1$ color-octet chiral super-multiplet $\hat{g}' = \{\tilde{g}', \sigma\}$ of extra gluinos and scalar σ fields to a vector hyper-multiplet of $N = 2$ supersymmetry. (Similarly, the electroweak sector, not to be discussed here, is supplemented by additional $SU(2)_L$ and $U(1)_Y$ super-multiplets.) The $N = 2$ mirror (s)fermions are assumed to be very heavy in order to avoid chirality problems. This hybrid model expands to $N = 2$ only in the gaugino sector [7].

2.1. The gluino sector

Standard MSSM $\tilde{g}$ and new gluinos $\tilde{g}'$ couple minimally to the gluon field

$$\mathcal{L}_{\text{SQCD}} \ni g_s \text{Tr} \left(\bar{\tilde{g}} \gamma^\mu [g_\mu, \tilde{g}] + \bar{\tilde{g}'} \gamma^\mu [g_\mu, \tilde{g}'] \right), \quad (1)$$

where g_s denotes the QCD coupling, the fields being color-octet matrices (*e.g.* $g_\mu = \frac{1}{\sqrt{2}}\lambda^a g_\mu^a$ with the Gell-Mann matrices λ^a); $\tilde{g}$ and $\tilde{g}'$ are two 4-component Majorana spinor fields. Quark and squark fields interact only with the standard gluino,

$$\mathcal{L}_{\text{SQCD}} \ni -g_s (\overline{q_L} \tilde{g} \tilde{q}_L - \overline{q_R} \tilde{g} \tilde{q}_R + \text{h.c.}) , \quad (2)$$

since only their hyper-multiplet partners (assumed to be heavy) couple to $\tilde{g}'$, as required by $N = 2$ supersymmetry.

Soft supersymmetry breaking generates masses for the gluino fields $\tilde{g}$ and $\tilde{g}'$,

$$\mathcal{M}_g = \begin{pmatrix} M'_3 & M_3^D \\ M_3^D & M_3 \end{pmatrix} . \quad (3)$$

Diagonal terms are induced by the individual Majorana mass parameters M_3 and M'_3 while an off-diagonal term corresponds to the Dirac mass. Diagonalization gives rise to two Majorana mass eigenstates, $\tilde{g}_1$ and $\tilde{g}_2$, with masses m_1 and m_2 . There are two limiting cases of interest: in the limit $M'_3 \rightarrow \pm\infty$ the standard MSSM gluino is recovered; in the limit of vanishing Majorana mass parameters M_3 and M'_3 with $M_3^D \neq 0$, the mixing is maximal and the two Majorana gluino states are paired to a Dirac state,

$$\tilde{g}_D = \tilde{g}_R + \tilde{g}'_L . \quad (4)$$

Dirac gluinos are characteristically different from Majorana gluinos; for detailed discussion we refer to [7]. Here we present one example from [7] in Fig. 1, where the partonic cross-sections for different-flavor squark production $qq' \rightarrow \tilde{q}\tilde{q}'$ mediated by the gluino t -channel exchange are plotted as a function of a Dirac/Majorana control parameter y , assuming partonic

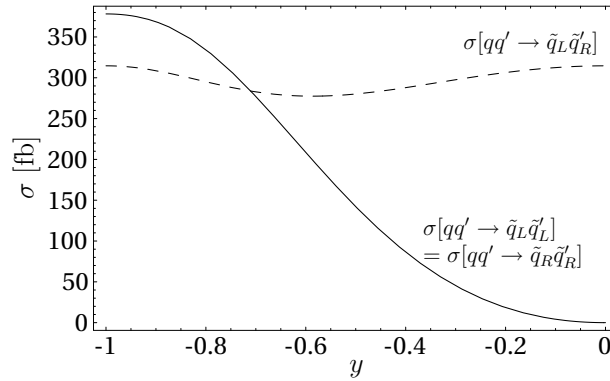


Fig. 1. Partonic cross-sections for different-flavor squark production as a function of the Dirac/Majorana control parameter y .

center-of-mass energy $\sqrt{s} = 2000$ GeV, $m_{\tilde{q}} = 500$ GeV and $m_{\tilde{g}_1} = 600$ GeV. Here $M'_3 = y M_3^D / (1 + y)$, $M_3 = -y M_3^D$ with $M_3^D = m_{\tilde{g}_1}$ kept fixed along $y \in [-1, 0]$. The parameter y allows for a continuous transition from $y = -1$, where the MSSM limit is reached with one Majorana gluino (the second being infinitely heavy), to $y = 0$ that corresponds to two degenerate Majorana fields combined to a Dirac gluino. For equal $\tilde{q}_L$ and $\tilde{q}_R$ masses, $\sigma(qq' \rightarrow \tilde{q}_L \tilde{q}'_L) = \sigma(qq' \rightarrow \tilde{q}_R \tilde{q}'_R)$ are non-zero in the Majorana limit but vanish in the Dirac, while $\sigma(qq' \rightarrow \tilde{q}_L \tilde{q}'_R)$ reaches the same value in both limits.

2.2. The sgluon sector

The new gluinos are accompanied by a color-octet complex scalar field σ . In the simplest realization, in parallel to the degenerate Majorana gluinos combined to the Dirac gluino, we assume that the real and imaginary components of the scalar field σ are degenerate with the (sgluon) mass denoted by M_σ [5]. At tree level, apart from the $\sigma\sigma^*g$ and $\sigma\sigma^*gg$ couplings as required by the gauge symmetry, the sgluons couple to the gluino pair via the Yukawa-type interaction

$$\mathcal{L}_{\tilde{g}_D \tilde{g}_D \sigma} = -\sqrt{2}i g_s f^{abc} \tilde{g}_{DL}^a \tilde{g}_{DR}^b \sigma^c + \text{h.c.}, \quad (5)$$

where f^{abc} are the $\text{SU}(3)_c$ structure constants. When supersymmetry is broken spontaneously, the Dirac gluino mass generates a scalar coupling between σ and squark pair [1]

$$\mathcal{L}_{\sigma \tilde{q} \tilde{q}} = -g_s M_3^D \sigma^a \frac{\lambda_{ij}^a}{\sqrt{2}} \sum_q (\tilde{q}_{Li}^* \tilde{q}_{Lj} - \tilde{q}_{Ri}^* \tilde{q}_{Rj}) + \text{h.c.}, \quad (6)$$

where L and R squarks contribute with opposite signs, as demanded by the general form of the super-QCD D-terms.

Although sgluons are R-parity even, the above couplings imply that tree-level single production of σ 's in gluon-gluon or quark-antiquark collisions is not possible. Only σ pairs can be produced in gluon collisions as well as in $q\bar{q}$ annihilation.

Since sgluons couple to squarks and gluinos, loop-induced couplings to the SM fields are of interest. However, even at the one-loop level gluino loops do not contribute to the σgg coupling as a consequence of Bose symmetry since the coupling is even in momentum space but odd, $\sim f^{abc}$, in color space¹. On the other hand, triangle diagrams involving squark/gluino lines

¹ Actually, the coupling of the octet sgluon to any number of gluons via the gluino loop is forbidden since the sgluon couples only to two different Majorana gluinos, Eq. (5), while gluons always couple to the same pair, Eq. (1).

(Fig. 2) generate σgg and $\sigma q\bar{q}$ couplings. The interaction Lagrangian however, Eq. (6), implies that all L- and R-squark contributions to the couplings come with opposite signs so that they cancel each other for mass degenerate squarks. In addition, the quark-antiquark coupling is suppressed by the quark mass as evident from general chirality rules.

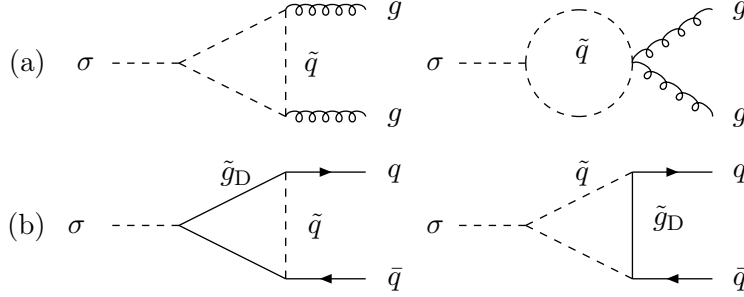


Fig. 2. Generic diagrams for the effective σgg (a) and $\sigma q\bar{q}$ (b) vertices with squark/gluino loops.

3. The phenomenology of sgluons at the LHC

Sgluons have a large color charge and thus might be copiously produced at the LHC. First we discuss their decay modes, and then production processes at the LHC.

3.1. Sgluon decays

Sgluons will decay into different channels that include gluinos, squarks, gluons and quarks. At tree level the σ particles can decay to a pair of Dirac gluinos $\tilde{g}_D$ or into a pair of squarks,

$$\Gamma[\sigma \rightarrow \tilde{g}_D \tilde{g}_D] = \frac{3\alpha_s M_\sigma}{4} \beta_{\tilde{g}} (1 + \beta_{\tilde{g}}^2), \quad (7)$$

$$\Gamma[\sigma \rightarrow \tilde{q}_a \tilde{q}_a^*] = \frac{\alpha_s |M_3^D|^2}{4M_\sigma} \beta_{\tilde{q}_a}, \quad (8)$$

where $\beta_{\tilde{g}, \tilde{q}_a}$ are the velocities of $\tilde{g}, \tilde{q}_a$ ($a = L, R$). If $M_\sigma < 2M_{\tilde{g}_D}, 2m_{\tilde{q}}$, one or both of these sparticles can be virtual. Squarks and gluinos will further cascade to SM particles and LSP.

Loop-induced couplings will mediate decays into gluon or quark–anti-quark pairs

$$\Gamma(\sigma \rightarrow gg) = \frac{5\alpha_s^3}{384\pi^2} \frac{|M_3^D|^2}{M_\sigma} \left| \sum_q [\tau_{\tilde{q}_L} f(\tau_{\tilde{q}_L}) - \tau_{\tilde{q}_R} f(\tau_{\tilde{q}_R})] \right|^2, \quad (9)$$

$$\Gamma(\sigma \rightarrow q\bar{q}) = \frac{9\alpha_s^3}{128\pi^2} \frac{|M_3^D|^2}{M_\sigma} m_q^2 \beta_q [(M_\sigma^2 - 4m_q^2) |\mathcal{I}_S|^2 + M_\sigma^2 |\mathcal{I}_P|^2]. \quad (10)$$

In Eq. (9), $\tau_{\tilde{q}_{L,R}} = 4m_{\tilde{q}_{L,R}}^2/M_\sigma^2$ and $f(\tau)$ is the standard function from a squark circulating in the loop [8]. In Eq. (10), the effective scalar (S) and pseudoscalar (P) couplings take the form ($\alpha = S, P$)

$$\mathcal{I}_\alpha = \int_0^1 dx \int_0^{1-x} dy [w_\alpha (C_L^{-1} - C_R^{-1}) + z_\alpha (D_L^{-1} - D_R^{-1})], \quad (11)$$

where $w_S = 1 - x - y$, $w_P = 1$, $z_S = (x + y)/9$, $z_P = 0$, and the squark/gluino denominators are ($a = L, R$)

$$C_a = (x + y) |M_3^D|^2 + (1 - x - y) m_{\tilde{q}_a}^2 - xy M_\sigma^2 - (x + y)(1 - x - y) m_q^2,$$

$$D_a = (1 - x - y) |M_3^D|^2 + (x + y) m_{\tilde{q}_a}^2 - xy M_\sigma^2 - (x + y)(1 - x - y) m_q^2.$$

If $\tilde{q}_L$ and $\tilde{q}_R$ of a given flavor mix, the subscripts L,R in equations above have to be replaced by 1, 2 labeling the squark mass eigenstates, and the contribution from this flavor is suppressed by the factor $\cos(2\theta_q)$. Note that, because of the chirality structure, decays to light quarks are suppressed by the quark mass and that both loop-induced decays are absent if L and R squarks are degenerate

The ordering between the tree-level and loop-induced decay modes depends on the values of various soft breaking parameters. So long as gluino-pair decay channels are shut kinematically, even for small L–R squark mass splitting the sgluon decays into two gluons, and to a $t\bar{t}$ pair if kinematically allowed, always dominate over tree-level off-shell four-body decays $\sigma \rightarrow \tilde{g}q\bar{q}\tilde{\chi}$ and $\sigma \rightarrow q\bar{q}\tilde{\chi}\tilde{\chi}$. Increasing the gluino mass increases the $\sigma\tilde{q}\tilde{q}^*$ coupling. As a result, the partial width into two gluons (due to pure squark loops) increases, while the $t\bar{t}$ partial width (due to mixed squark–gluino loops) decreases since the increase of the $\sigma\tilde{q}\tilde{q}$ couplings is over-compensated by the gluino mass dependence of the propagators. Of course, the tree-level, two-body decays of Eqs. (7),(8) will dominate if they are kinematically allowed. Well above all thresholds the partial width into gluinos always dominates:

it grows $\propto M_\sigma$ while the partial width into squarks asymptotically scales like $1/M_\sigma$ since the supersymmetry breaking $\sigma\tilde{q}\tilde{q}^*$ coupling has mass dimension 1, while the supersymmetric $\sigma\tilde{g}\tilde{g}$ coupling is dimensionless.

The above qualitative features can be seen in Fig. 3, where the branching ratios for σ decays are plotted for two different squark/gluino mass hierarchies. Moderate mass splitting between the L and R squarks of the five light flavors, and somewhat greater for soft breaking t masses have been assumed: $m_{\tilde{q}_R} = 0.95m_{\tilde{q}_L}$, $m_{\tilde{t}_L} = 0.9m_{\tilde{q}_L}$, $m_{\tilde{t}_R} = 0.8m_{\tilde{q}_L}$, the off-diagonal element of the squared $\tilde{t}$ mass matrix $X_t = m_{\tilde{q}_L}$; and the gluino is a pure Dirac state, *i.e.* $m_{\tilde{g}} = |M_3^D|$.

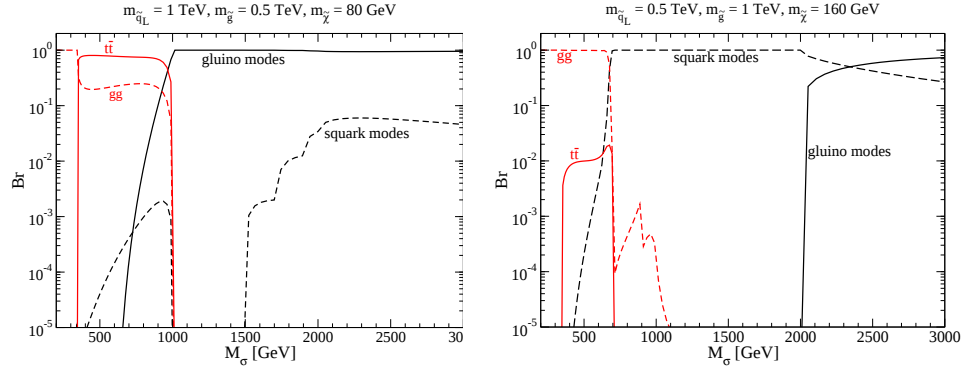


Fig. 3. Branching ratios for σ decays, for $m_{\tilde{q}_L} = 2m_{\tilde{g}} = 1$ TeV (left) and $m_{\tilde{g}} = 2m_{\tilde{q}_L} = 1$ TeV (right).

3.2. Sgluon production and signatures at the LHC

Sgluons can be produced singly in gluon–gluon collisions via squark loops; the production via $q\bar{q}$ annihilation is negligible for light incoming quarks. With the Breit–Wigner coefficient split off, the partonic cross-section is given by

$$\hat{\sigma}[gg \rightarrow \sigma] = \frac{\pi^2}{M_\sigma^3} \Gamma(\sigma \rightarrow gg), \quad (12)$$

where the partial width for $\sigma \rightarrow gg$ decays is given in Eq. (9). The resulting cross-section for single σ production at the LHC is shown by the blue (c) and (d) curves in Fig. 4 (with the LO CTEQ6L parton densities [9]). The solid curve (c) has been calculated for the parameter set of the right frame of Fig. 3, while the dashed one (d) is for the benchmark point SPS1a' [10]. Since SPS1a' has a somewhat smaller gluino mass (interpreted here as a Dirac mass) it generally leads to smaller cross-sections for single σ production.

The signatures for single σ production are potentially very exciting. However, the 2-gluon decay channel must be discriminated from the large SM background. On the other hand, large production cross-section in the gluon fusion would imply diminished decay rates to other channels, which in addition do not allow a direct reconstruction of M_σ . Detailed experimental simulations are needed to see if the single σ production can be detected as a resonance above the SM and MSSM backgrounds.

Sgluons can also be pair-produced in $q\bar{q}$ and gg processes,

$$\sigma[q\bar{q} \rightarrow \sigma\sigma^*] = \frac{4\pi\alpha_s^2}{9s} \beta_\sigma^3, \quad (13)$$

$$\sigma[gg \rightarrow \sigma\sigma^*] = \frac{15\pi\alpha_s^2\beta_\sigma}{8s} \left[1 + \frac{34}{5} \frac{M_\sigma^2}{s} - \frac{24}{5} \left(1 - \frac{M_\sigma^2}{s} \right) \frac{M_\sigma^2}{s} L_\sigma \right], \quad (14)$$

where $\sqrt{s}$ is the invariant parton-parton energy, M_σ and β_σ are the mass and center-of-mass velocity of the σ particle, and $L_\sigma = \beta_\sigma^{-1} \log(1 + \beta_\sigma)/(1 - \beta_\sigma)$. The cross-section for σ -pair production at LHC, $pp \rightarrow \sigma\sigma^*$, is shown by the solid red curve (a) in Fig. 4 for the σ -mass range between 500 GeV and 2 TeV. With values from several picobarn downwards, a sizable $\sigma\sigma^*$ event rate can be generated. As expected, due to large color charge of the sgluon,

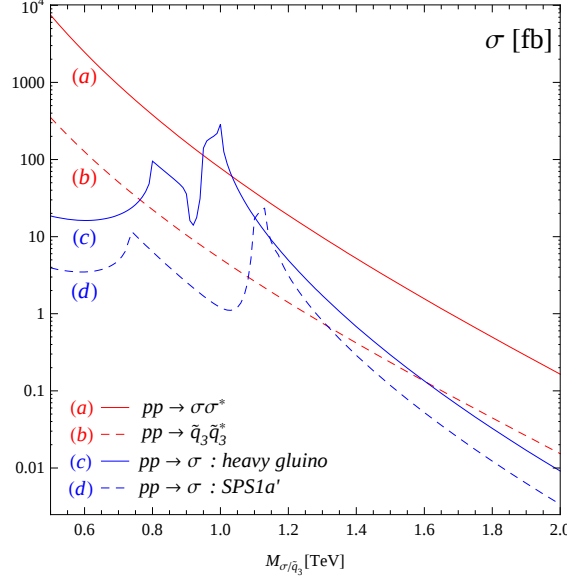


Fig. 4. Cross-sections for σ -pair (and $\tilde{q}_3$ -pair) production (red lines (a) and (b)), as well as for single σ production (blue lines (c) and (d)), at the LHC. The solid blue curve (c) is for the mass parameters as in Fig. 3 (right), while the dashed blue (d) for the mSUGRA benchmark point SPS1a'.

the $\sigma\sigma^*$ cross-section exceeds stop or sbottom-pair production (red dashed line (b)), mediated by a set of topologically equivalent Feynman diagrams, by more than an order of magnitude. In the scenario of the right frame of Fig. 3, the single σ cross-section can exceed the σ -pair production cross-section for $M_\sigma \sim 1$ TeV; taking $m_{\tilde{q}} \simeq 2|M_3^D|$, as in the left frame of Fig. 3, lead to a very small single σ production cross-section. In general one cannot simultaneously have a large $\sigma(pp \rightarrow \sigma)$ and a large $\text{Br}(\sigma \rightarrow t\bar{t})$.

With the exception of $\sigma \rightarrow gg$ decays, all the σ decay modes give rise to signatures that should be easily detectable. Most spectacular signature results from $\sigma \rightarrow \tilde{g}\tilde{g}$ decay, each σ decaying into at least four hard jets and two invisible neutralinos as LSP's. σ -pair production then generates final states with a minimum of eight jets with high sphericity and four LSP's. In such events the transverse momenta of the hard jets produced and the vector sum of the momenta of the four χ_1^0 in the final state, which determines the measured missing transverse momentum p_T , are markedly different from the corresponding MSSM gluino or squark production with the same mass configurations [5].

Other interesting final states are four-stop states $\tilde{t}_1\tilde{t}_1\tilde{t}_1^*\tilde{t}_1^*$, which can dominate if $m_{\tilde{q}} \lesssim m_{\tilde{g}}$ and L-R mixing is significant in the stop sector, and $\tilde{q}\tilde{q}^*\tilde{g}\tilde{g}$ if $M_\sigma > 2m_{\tilde{g}} \gtrsim 2m_{\tilde{q}}$. These channels also lead to four LSPs in the final state, plus a large number of hard jets. On the other hand, the $t\bar{t}\bar{t}\bar{t}$ final state, which can be the dominant mode if the two-body decays into squarks and gluinos are kinematically excluded, might allow the direct kinematic reconstruction of M_σ . In addition, the observation of $t\bar{c}\bar{t}c$ final states might indicate a substantial mixing in the up-type squark sector.

4. Summary

Models with Dirac gauginos offer an interesting alternative to the MSSM scenario. Embedded into theories of extended supersymmetries, they predict the presence of scalar particles in the adjoint representation of the gauge groups. The color-octet scalars, sgluons, (if kinematically accessible) can be copiously produced at the LHC. Their signatures are distinctly different from the usual MSSM topologies. Depending on the masses of the particles involved, either multi-jet final states with high sphericity and large missing transverse momentum are predicted, or four top quarks should be observed in 2σ production. If the mass splitting between L and R squarks is not too small, loop-induced single σ production may also have a sizable cross-section. Though this channel suffers in general from large backgrounds, identifying the σ particle as a resonance in 2-gluon final states would truly be an exciting experimental observation.

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