Seeking Sgluons

Tilman Plehn

SUPA, School of Physics and Astronomy, University of Edinburgh, Scotland

Tim M.P. Tait

Argonne National Laboratory, Argonne, IL 60439
Northwestern University, 2145 Sheridan Road, Evanston, IL 60208

Scalar gluons — or sgluons — are color octet scalars without electroweak charges. They occur in supersymmetric models of Dirac gauginos as the scalar partners of the gluino and carry Standard-Model type $R$ charge. This allows them to interact with ordinary matter and to be produced at the LHC, singly as well as in pairs. Sgluons dominantly decay into gluons, top pairs, and a top quark plus a light quark. A pair of sgluons decaying into like-sign tops would provide a striking signature at the LHC. In our discussion of this channel we especially focus on the proper treatment of QCD jets.

I. INTRODUCTION

Theories with weak scale supersymmetry represent the most complete and cherished vision of physics beyond the Standard Model. Their many successes include stabilization of the electroweak scale with respect to high-scale physics, improvement of the convergence of couplings necessary for Grand Unification, possible electroweak baryogenesis to explain the matter-anti-matter asymmetry of the Universe, and (with $R$ parity) relatively mild contributions to precision electroweak data and a successful dark matter candidate.

However, none of those successes rest crucially on the minimal realization of the supersymmetric Standard Model [1]. In fact, there are features of the minimal supersymmetric standard model (MSSM) which are somewhat at odds with the MSSM as a completely natural theory of the electroweak scale. The $\mu$ or $B\mu$ problem indicates that the supersymmetric $\mu$ term must be roughly of the same size as the supersymmetry-breaking parameters, and yet the MSSM offers no explanation for why this should be the case. The lightest Higgs mass in the MSSM is, at tree-level, less than the $Z$ boson mass, and even including radiative corrections typically in conflict with the LEP-II bound [2]. The model survives based on large radiative corrections from scalar top quarks, but in turn this requires the stops to be so heavy that the Higgs soft masses end up fine-tuned to the per-cent level.

Perhaps the most disturbing feature of the MSSM is the fact that if one naively assumes a general spectrum of supersymmetry-breaking parameters, it is simply ruled out. If there are large mixings among the squarks and among the sleptons, flavor-violating processes such as $K\bar{K}$ mixing, $\mu \rightarrow e\gamma$, and others, can be enhanced by orders of magnitude with respect to Standard Model predictions, in obvious contradiction to experimental data. The traditional solution (of which gauge mediation is the prototype) is to engineer supersymmetry breaking such that the soft-breaking mass parameters are highly flavor diagonal [3]. This controls flavor violation to an acceptable level at a given energy scale, but continues to be challenged by natural electroweak symmetry breaking, the Higgs mass, and cosmology.

An alternative is to alter the low energy structure of the model such that the supersymmetric solution to the hierarchy problem is preserved, but flavor violation is ameliorated. A recent elegant solution imposes a continuous $R$ symmetry [4], to form the minimal $R$-symmetric supersymmetric Standard Model (MRSSM) [1]. This model invokes the $R$ symmetry to forbid the potentially flavor-violating left-right mixing $A$-terms, and to guarantee (at least largely in light of anomaly-mediated contributions) Dirac gaugino masses. This Dirac nature often leads to an additional suppression in flavor-violating processes from gauginos running in the loops, which causes many contributions to scale with large gaugino mass as $1/m_\tilde{g}^2$ instead of the Majorana scaling, $1/m_\tilde{g}$ [1]. The model has additional interesting features, such as UV finite scalar masses due to the super-soft feature of Dirac gauginos [6], and additional matter content at the electroweak scale. For example, if the gluino is a Dirac particle, the two on-shell degrees of freedom from the gluon are not sufficient to construct a supersymmetric four-spinor. To provide the two additional bosonic degrees of freedom, the MRSSM adds a complex scalar partner. An unbroken $R$ symmetry is also a generic feature of simple models of meta-stable supersymmetry breaking [7].
In this article we consider the phenomenology of the scalar partner of the Dirac gluino, a color adjoint scalar — the sgluon. As a colored particle, it couples to a gluon and will have large pair-production rates at the Tevatron and LHC. It interacts with quarks at the one-loop level proportional to the quark masses. Thanks to the large mixing naturally expected in the MRSSM’s squark sector, it will readily decay into flavor violating channels. Previous studies have considered scalar color octets with either decays into missing energy [8], with electroweak as well as SU(3) charges [9], or purely flavor diagonal couplings [10], which all lead to very different phenomena.

Our work is organized as follows. In Section II we show how sgluons arise in a model with Dirac gluinos, examine their soft masses, and derive their interactions including the one-loop contribution to the $G-q-q'$ and $G-g-g$ vertices. In Section III we examine the flavor-violating effects mediated by sgluons and derive some mild constraints from $K$-$\bar{K}$ mixing. In Section IV we discuss the production of a pair of sgluons through the strong interaction, and examine the decay of the pair into like-sign tops. This distinctive signature provides bounds on the sgluon mass at the Tevatron and will easily be discovered at the LHC for a wide range of masses. We particularly focus on the proper treatment of the jet activity in such events.

II. SGLUONS

Scalar gluons are contained in a chiral superfield $\Phi^a$ which is a color adjoint carrying $R$ charge zero. The fermionic component $\psi^a$ is married through $D$–term supersymmetry breaking to the ordinary gluino $\lambda^a$. The lowest component $G^a$ is a complex color adjoint scalar. Supersymmetry breaking will generally split this into two real scalar states which are admixtures of the real and imaginary parts of $G$. We discuss the spectrum below.

Kinetic terms for the sgluons are contained in canonical Kähler potential terms for $\Phi$,

$$ \int d^4\theta \; \Phi^\dagger e^{-V} \Phi $$

where $V$ is the vector superfield containing the gluon and the $SU(3)_C$ gauge indices are implied. The kinetic terms include the coupling of the sgluon $G$ to the gluons from the covariant derivative, a $G$-$\psi$-$\lambda$ coupling of strength $g_s$ required by supersymmetry, and $D$–term contributions to the scalar potential that are of the form $G^*G-\bar{q}^\gamma q$ which will not be important for our purposes. There are no renormalizable gauge invariant terms through which $\Phi$ interacts with matter superfields in either the Kähler potential or in the super-potential, and the assumed $R$ symmetry is incompatible with $\Phi^2$ or $\Phi^3$ interactions in the super-potential. Thus, the tree-level supersymmetric interactions of $G$ are determined entirely by supersymmetric QCD,

$$ \mathcal{L}_{\text{SQCD}} = (D_{\mu} G)^* (D^{\mu} G) + i\sqrt{2} \ g_s f_{abc} \bar{\tilde{g}}^b (G^a P_L + G^{*a} P_R) \tilde{g}^c $$

where $D^\mu$ is the usual covariant derivative for a color adjoint, $\tilde{g}$ is the (four-component) gluino, and $f_{abc}$ are the structure constants of $SU(3)$.

A. Supersymmetry breaking and masses

Soft mass terms for the sgluons can arise from either $F$–term ($\langle X \rangle = \theta^2 F$) or $D$–term ($\langle W' \rangle = \theta D'$) spurions of supersymmetry breaking,

$$ \int d^4\theta \left\{ \frac{1}{M_1^2} X^\dagger X \Phi^\dagger \Phi + \frac{1}{M_2^2} X^\dagger X \text{Tr} \Phi^2 \right\} + \int d^2\theta \frac{1}{M_3^2} W'^\alpha W^{*\alpha} \text{Tr} \Phi^2 + \text{H.c.} $$

and also are generated by the term responsible for the Dirac gluino mass [8],

$$ \int d^2\theta \frac{\sqrt{2}}{M_4} W^{*\alpha} W_{3\alpha} \Phi^a $$
where \( W_{3\mu}^a \) is the usual superfield \( SU(3)_C \) field strength. This super-potential term, along with the usual MSSM \( SU(3)_C \) \( D \)-terms, lead to terms in the Lagrangian,

\[
-m_\tilde{g} \lambda^a \psi^a - \sqrt{2} \left( m_\tilde{g} G^a + m_\tilde{g}^* G^{a*} \right) D^a - g_s D^a \sum_{\tilde{q}_L} \tilde{q}_L^a T^a \tilde{q}_L + g_s D^a \sum_{\tilde{q}_R} \tilde{q}_R^a T^a \tilde{q}_R - \frac{1}{2} D^a D^a 
\]

(5)

where \( m_\tilde{g} = D'/M_4 \) is the Dirac gluino mass, and \( T^a \) are the generators of \( SU(3)_C \) in the fundamental representation. Replacing the \( SU(3) \) auxiliary field \( D^a \) through its equation of motion leads to terms proportional to \( m_\tilde{g}^2 G^2 \) and \( m_\tilde{g}^2 G^{a*2} \) as well as \( |m_\tilde{g}| G^2 \). It also induces tri-linear interactions of \( G \) with squarks, somewhat analogous to the \( A \) terms in the usual MSSM. Altogether, the supersymmetry-breaking Lagrangian for \( G \) reads

\[
\mathcal{L}_{\text{soft}} = m_1^2 |G^a|^2 + \frac{1}{2} m_2^2 G^{a2} + \frac{1}{2} m_2^2 G^{a*2} - \sqrt{2} g_s \left( m_\tilde{g} G^a + m_\tilde{g}^* G^{a*} \right) \left( \sum_{\tilde{q}_L} \tilde{q}_L^a T^a \tilde{q}_L - \sum_{\tilde{q}_R} \tilde{q}_R^a T^a \tilde{q}_R \right) 
\]

(6)

where \( m_1^2 \) is a real parameter and \( m_2^2 \) may be complex. The mass eigenstates are two real color adjoint scalars which can be labelled as \( G_1 \) and \( G_2 \). The mass-squared eigenvalues are given by,

\[
m^2_{G_1,G_2} = m_1^2 \mp |m_2^2| ,
\]

(7)

and clearly we must have \( m_1^2 > |m_2^2| \) or run the risk of a color-breaking vacuum. There will be a non-trivial mixing angle when \( m_2^2 \) is complex. When we write \( m_2^2 \) in terms of its phase \( m_2^2 = |m_2^2|e^{i\gamma} \), the mass eigenstates are

\[
G_1^a = \sin \frac{\gamma}{2} G^a + \cos \frac{\gamma}{2} G^{a*} , \\
G_2^a = \cos \frac{\gamma}{2} G^a - \sin \frac{\gamma}{2} G^{a*} .
\]

(8)

For simplicity we will assume \( m_\tilde{g} \) and \( m_2^2 \) are real, and thus there is no non-trivial mixing from here on. In that case both sgluons are either a pure scalar or a pure pseudoscalar, which is equivalent as long as we combine them with massless QCD, the theory relevant for LHC.

All tree-level interactions of the sgluon with Standard Model and MSSM states we can read off \( \mathcal{L}_{\text{SQCD}} \) and \( \mathcal{L}_{\text{soft}} \). The coupling of two sgluons to gluons is simply a result of its adjoint color charge and arises from the kinetic term. The coupling to two gluinos is the supersymmetric partner of the gluon couplings, while the couplings to two squarks arise from \( D \) terms. Note in particular that the Dirac gluino mass sets the size of the squark-squark-sgluon coupling.

### B. Loop-induced coupling to quarks

If we insert the squark and gluino couplings to a sgluon shown in eq. (2) and in eq. (3) into one-loop diagrams, the sgluon will couple to quarks. There are two Feynman graphs responsible for this interaction induced by gluinos and squarks, shown in Figure 1. For each of the two sgluon states the effective action after electroweak symmetry breaking contains a dimension-4 operator of the form,

\[
G^a \left[ \tilde{q}^i \, T^a \left( g_L^i P_L + g_R^i P_R \right) \, q^j \right] + \text{H.c.}
\]

(9)
where $i$ and $j$ are quark flavor indices. The couplings may be expressed as

$$g_{ij}^{(L/R)} = \frac{g_3 \delta_{ij}}{16\pi^2} \frac{m_{\tilde{q}}}{m_G^2} \left( m_{i,f}^{(L/R)} - m_{j,f}^{(L/R)} \right)$$

where $\delta_{ij}$ is the relevant squark mixing parameter, $m_{i,j}$ are the quark masses, and the dimensionless $f_{i,j}$ are functions of the heavy sgluon, gluino, and the squark masses. Their form is given in the Appendix. It is important to notice that in the MRSSM these couplings come out automatically proportional to quark masses, which mitigates their contribution to flavor-violating observables. In contrast, for example, the $A$ terms in the MSSM have to be defined to appear proportionally with the quark masses, i.e. as $m_q A_q$, which is an additional assumption on the flavor structure of the general MSSM. In the limit of degenerate squarks, this source of flavor violation in the MRSSM will switch off through a super-GIM mechanism.

In the limit of large squark, gluino, or sgluon masses the quark-quark-sgluon coupling will be suppressed by $m_{\tilde{q}}/m_{\tilde{G}}^2$, $1/m_{\tilde{q}}$, or $m_{\tilde{g}}/m_{\tilde{G}}^2$, respectively.

C. Sgluon-gluon-gluon Coupling

Pairs of sgluons interact with one or two gluons as a consequence of the fact that the sgluons are adjoints of $SU(3)$. There is also a single sgluon-gluon-gluon interaction mediated by squarks, as shown in Figure 2. The effective action contains a dimension-five operator generated at one loop

$$\frac{g_3^2}{16\pi^2} \frac{m_{\tilde{q}}}{m_G^2} d^{abc} \lambda_g G^{a}_{\mu\nu} F_{bc}^{\mu\nu} + \text{H.c.}$$

where $d^{abc}$ is the symmetric gauge-invariant combination of three $SU(3)$ adjoints and the dimensionless form factor $\lambda_g$ is given in the Appendix. A similar contribution mediated by gluinos vanishes at one loop, due to the symmetry structure of the (asymmetric) color factor and the (symmetric) loop contribution. In the limit of heavy squarks the coupling $\lambda_g$ will vanish proportionally to $1/m_{\tilde{q}}^2$.

III. SGLUON-MEDIATED FLAVOR VIOLATION

Through its flavor-violating couplings to quarks, the sgluon can also mediate flavor-changing processes. At energies far below the sgluon mass, these interactions look like $\Delta F = 2$ four fermion interactions. Using the $s$-$d$ transition operators relevant for $K-\bar{K}$ mixing we can illustrate the terms left behind when the sgluon is integrated out:

$$- \frac{1}{2m_G^2} \left\{ (g_L^{sd} + 2g_R^{sd}) \left[ Q_5^{sd} - \frac{1}{3} Q_4^{sd} \right] + g_R^{sd} g_L^{sd} \left[ Q_3^{sd} + \bar{Q}_3^{sd} - \frac{1}{3} Q_2^{sd} - \frac{1}{3} \bar{Q}_2^{sd} \right] \right\}$$

where the $Q_i^{ij}$ are the four fermion interactions defined in Ref. [11]. Comparing to eq. (10) we know that the couplings $g$ are proportional to $\delta$, so the relevant parameters which will be constrained are $\delta/m_G$. Analogous expressions describe $b$-$d$, $b$-$s$, and $c$-$u$ mixing. We compute the coefficients of each of these
operators and compare with the global fit by the UTfit collaboration [11] to obtain bounds on the sgluon mass, for a given choice of gluino mass, average squark masses, and squark flavor mixing parameters $\tilde{\delta}$. In Fig. 3, we present the scale of the most constraining of the four fermion operators, expressed as the effective scale $\Lambda / \tilde{\delta}^{ij}$. Since mixing effects will only raise the effective scale of the four fermion interactions, the curves are the minimum possible effective scales for a given choice of sgluon, gluino, and average squark masses. In other words, a realistic choice of mixing parameters will increase the effective scales which suppress the FCNC operators and thus result in less constraints from flavor violation mediated by sgluons. We also show the bounds on the corresponding scales from Ref. [11]. Bounds from $B-\bar{B}$ and $D-\bar{D}$ mixing are mild enough as to basically provide no constraint. The $K-\bar{K}$ mixing bound is more severe, but the plotted bound is on the imaginary part of $\tilde{\delta}$, and thus can be avoided if there is no large $CP$-violating phase in the down-strange squark mixing element. In any case, sgluon masses above 600 GeV or so are compatible with measurement, regardless of mixing. Note also that the constraints on down-type mixing do not in any way preclude the sgluon flavor-violating decays into up-type quarks which we consider below.

IV. SGLOUNS AT COLLIDERS

As discussed in Section [11] we expect the $R$-symmetric supersymmetric theory to include a pair of sgluons whose mass is split by something of order the gluino mass. There will thus be two sgluon states which we can search for at hadron colliders. Since the masses are not typically degenerate, the results in this section are presented for a single sgluon state, and apply equally to the lighter or the heavier sgluon. The couplings to quarks and single sgluon coupling to gluons will depend on the mixing between the two states, but tree level pair production of sgluons involves only the strong coupling, as these interactions are protected by $SU(3)_C$ gauge invariance. The resulting pair cross sections thus only depend on the sgluon mass, similar to, for example, the case of scalar leptoquark pairs.

In Figure 4, we present the leading order cross section for pair production of sgluons as a function of their mass at the Tevatron and LHC [12]. For masses around 250 GeV, the production rate of the order of 200 fb at the Tevatron would correspond to a few hundred sgluon pair events in the currently available CDF and DZero data, depending on the decays and triggers. As expected, the production rate
FIG. 4: Inclusive production cross sections for sgluons at the Tevatron and at the LHC. For the LHC we show pair production (solid) and single production (dashed) as a function of the sgluon mass. The two curves for single sgluon production assume a gluino mass of 1 TeV and squark masses of 500 GeV (upper curve) and 1 TeV (lower curve).

at the Tevatron drops below the femtobarn level for sgluon masses around 400 GeV, thanks to the limited center-of-mass energy and $p$-wave suppression of the dominant subprocess $q\bar{q}\rightarrow GG$. From a dedicated analysis we could expect sgluon mass bounds similar to squark mass bounds in the limit of large gluino mass.

At the LHC, the dominant subprocess is $gg\rightarrow GG^*$ with large cross sections, falling from around 400 pb for masses around 250 GeV to 2 fb for masses around 1.5 TeV. For sgluon masses in the TeV range, the LHC will rely on its sea-quark suppressed $q\bar{q}$ luminosity which leads to a rapid drop of the cross section above 1.8 TeV. With an appreciable LHC luminosity these rates correspond to several hundred to a few million events available for analyses.

For the LHC we also present the single production rate from gluon fusion \[\text{[13]}\], for two choices of squark and gluino masses. Through the one-loop diagrams discussed in Sec. [1] there can be appreciable single sgluon production through $gg\rightarrow G$, which can dominate for large sgluon masses because of the phase space suppression of the pair production and threshold effects. For small sgluon masses the LHC is not energy limited, so the single production channel is suppressed by a loop factor $\alpha_s/(4\pi)$ squared. The problem of single production will be challenging backgrounds discussed below. The sgluon interaction with quarks typically results in a negligible single sgluon production rate, because the leading contribution is one loop and the suppression by the light quark masses.

Sgluons can also be produced at the LHC in cascade decays involving squarks, through the soft breaking interaction of eq.(6). Production of pairs of heavier squarks can thus cascade down through sgluons into lighter squarks, leaving behind either light jets which reconstruct the sgluon mass, or events enriched with top quarks. Either of these possibilities is rather exotic from the point of view of the standard MSSM, and we leave their detailed exploration for future work.

A. Sgluon decays — like-sign tops

If heavy enough, sgluons will decay at tree level into pairs of gluinos and/or squarks. If these decay channels are closed, the sgluon has to decay through its loop-induced couplings into quarks and gluons. In the case of decays into quarks, the $G-q\bar{q}$ interaction is proportional to the heavier of the two quark masses, so one can expect that decays including at least one top quark to dominate.

If the mixing in the up-type squark sector is large (as is the point of the MRSSM), we therefore expect large and comparable branching ratios into $\bar{t}u$, $\bar{t}c$, and $\bar{t}t$. The specifics of the branching ratios are a window into the details of the squark mixing matrices. In Figure [5] we present the branching ratios as
FIG. 5: Branching ratios for sgluon decays into $gg$, $t\bar{t}$, $qt + \bar{t}q$ for $q = u, c$ and $\tilde{q}\tilde{q}^*$ as a function of sgluon mass and for two choices of left-handed squark and gluino masses. Right-handed squark masses are set to 90% of the left-handed squark masses. We assume maximal up-squark mixing.

a function of the sgluon mass for two sets of average squark and gluino masses, and assuming maximal mixing in the up-squark sector. The mixed heavy–light quark decays can dominate the sgluon decays for small masses and decreases with larger sgluon masses. For the maximal mixing considered here, the decay into $t\bar{t}$ is roughly comparable to any single one of the heavy-light decays. Thus, combing channels together, the branching ratio into one top (or anti-top) and a light quark are a factor of a few times larger than that into a top pair. For heavier squarks the supersymmetric decay channels are typically closed in the region accessible to the LHC. Again, the decay to one heavy and one light quark typically dominates. At large sgluon masses decays into gluons dominate because the $G–q–g$ coupling is a dimension five operator which grows with the invariant mass of the sgluon, i.e. the only scale in the process, while the decay to quarks will be suppressed by a relative factor $m_t/m_G$.

When sgluons are pair-produced, each sgluon is as likely to decay into a top as an anti-top. Thus, half of the decays where both sgluons decay into a top and a light-quark initiated jet will have same sign tops ($tt$ or $\bar{t}\bar{t}$). When both tops decay leptonically we have a final state containing two light jets and either $b\ell^+\nu$ $\bar{b}\ell^+\nu$ or a $b\ell^-\bar{\nu}$ $\bar{b}\ell^-\bar{\nu}$ (where $\ell$ is an electron or muon) - a striking signature of physics beyond the Standard Model that also arises in the context of models of top compositeness [14].

Reference [15] has considered a search for like-sign tops at the Tevatron in the context of a model of maximal flavor violation. The authors perform a sophisticated treatment of the backgrounds and the CDF detector efficiencies. While the maximal flavor violation model signal is the result of a mixture of pair and single production of a neutral color singlet scalar $\eta$ which interacts moderately strongly with top and charm, the analysis only requires a pair of like sign leptons, a $b$-tagged jet, and missing energy. Thus, our signal events are expected to have a high efficiency with respect to the analysis cuts, and one could get a Tevatron bound on the sgluon mass from a similar analysis.

A single sgluon produced at the LHC can decay through its flavor-violating interactions into a single top quark and a light jet, similar in topology to the $s$-channel mode of single top production. Even in the Standard model, this mode is challenging at the LHC because of large backgrounds from $t\bar{t}$ and $t$-channel single top [16]. Therefore, single sgluon production is unlikely to be phenomenologically relevant after considering QCD effects and backgrounds. However, if the sgluon is heavy enough it may be possible to use the peak in the top plus light jet invariant mass to isolate a signal [17].
B. LHC signatures

From the discussion above we are immediately lead to study in more detail sgluon pair production with a subsequent decay into two like-sign tops with leptonic decays. With sizeable production rates and branching ratios above 10% the question remains if this channel survives the LHC triggers and acceptance cuts. In Fig. 6 we show the normalized transverse momentum distributions of all sgluon decay particles. For both sgluon masses of 300 and 600 GeV the light-flavor jet and bottom transverse momenta peak above $p_T > 50$ GeV, which indicates that observing four jets and tagging two bottoms should not be a problem. For a 300 GeV sgluon the harder jet has a typical transverse momentum around $p_{T,j} \sim m_G - m_t \sim 120$ GeV. Similarly, the harder bottom can acquire $p_{T,b} \sim m_t - m_W \sim 100$ GeV. For a heavier sgluon we see that the bottom distributions hardly change, since they are mostly determined by the top and $W$ masses. The two light-flavor jets from the sgluon decay become significantly harder and peak around $p_{T,j} \sim 180$ GeV and 280 GeV, respectively. As we will see later, in particular for heavier sgluons the decay jet can be identified unambiguously even in the presence of QCD jets.

On the lepton side of Fig. 6 we see that the harder lepton with a typical transverse momentum close to 100 GeV guarantees the triggering of sgluon pair production. Moreover, such a large transverse momentum might be useful to distinguish the sgluon pairs from the Standard Model backgrounds, even in the same-sign lepton case. The second lepton is comparably soft, and the missing transverse momentum peaking around 70 GeV is unlikely to contribute to the smoking-gun signature.

Nevertheless, from the distributions in Fig. 6 we can see that triggering and acceptance cuts will not be a problem for the largely background-free like-sign tops signature. Moreover, both hard light-flavor decay jets as well as the relatively hard lepton spectrum should help to reconstruct the sgluon mass scale, which would allow us to gain information on the branching ratio and thereby on the flavor structure of the MRSSM.

C. Sgluons and QCD jets

Because the light-flavor decay jets from the sgluons are crucial for the analysis described above, we have to answer the question if this jet can be identified in the jet-rich LHC environment. This question becomes even more relevant once we try to search for the opposite-sign lepton channel or attempt to determine...
FIG. 7: Number of jets produced in QCD jet radiation in addition to the sgluon or top decay jets. We use MLM merging as implemented in MadEvent and apply $p_{T,j} > 30, 50, 100$ GeV and $\Delta R_{jj} > 0.4$ on the parton level.

We simulate additional QCD jets in the signal process $pp \to GG + X$ using the MadEvent implementation of the MLM scheme. In particular because the sgluons in the final state provide a hard scale for the process we do not expect the MLM results to differ from a corresponding CKKW analysis. In particular for heavier sgluon masses the QCD activity should be dominated by collinear parton-shower effects, but the MLM scheme now allows us to consistently treat top backgrounds and the sgluon signal for different masses. For this simulation we avoid introducing a supersymmetric shower including sgluon splittings, which would be required to include final-state radiation. Because of the lack of a collinear enhancement we know that final-state radiation will not contribute strongly to QCD jet radiation. Therefore, we only include initial-state radiation which is universal for different heavy new-physics states produced at the LHC.

In Fig. 7 we show the number of QCD jets in sgluon pair events at the LHC. Apart from a crucial minimal jet separation of $R_{jj} > 0.4$ we apply only a varying transverse momentum cut of 30, 50, and 100 GeV on the radiated jets. In the left panel we see that while top quarks most likely come with no additional jet from initial-state radiation, a 300 GeV sgluon will most likely be accompanied by one and a 600 GeV sgluon by two additional QCD jets. This is an effect of the hard factorization scale in the process which determines the size and the maximum range of the collinear enhancement of initial-state radiation. When we increase the minimum $p_{T,j}$ to 50 GeV the typical number of additional jets drops by roughly one, but in particularly heavy states still come with zero, one, or two jets at roughly the same rate. Only the three additional jets channel is suppressed to the 10% level, where this quantitative result should be taken with a grain of salt without a tuned parton shower for the LHC. Finally, a cut of at least $p_{T,j} > 100$ GeV gets rid of additional jets in roughly two thirds of the events and allows us to use light-flavor decay jets in the analysis.

The normalized transverse momentum distributions for the radiated jets are shown in Fig. 8. The different areas under the curves for the four leading jets reflect the fact that only a fraction of events actually show such jets. Again, we see that from the top pairs to the 600 GeV sgluons the situation changes qualitatively: in the bottom panel the curves for the leading and the sub-leading are similar, and as long as we stay below $p_{T,j}$ even a third QCD jet is very likely to appear. In a way, the crossing point between the two hardest jet indicates a $p_{T,j}$ range below which we should not rely on QCD jets being rare or suppressed. In other words, 50 GeV jets should not be used as part of the signal unless we have a way to identify decay jets, while jets with $p_{T,j} > 100$ GeV are comparably safe. Of course, this insight is not new — for example, careful squark and gluino analyses for the LHC have always been designed that way.

The good news of this QCD analysis is that jet activity will actually be helpful to distinguish sgluon pairs from $t\bar{t}$ backgrounds. First, top pair events are quite unlikely to come with a QCD jet of $p_{T,j} >
100 GeV, which we have seen in Sec. [V.B] is typical for the signal's decay jets. Secondly, as long as we require $p_{T,j} > 100$ GeV for the decay jets, even a 600 GeV sgluon will hardly come with such hard QCD jets. In general, if we consider something like the total visible mass as an observable to distinguish the sgluon-pair signal from top pairs, Fig. [7] shows that QCD radiation will improve this handle. Based on these results, we can estimate that for sgluon branching ratios into $t\bar{q}$ of order 10%, we can expect an LHC discovery reach of up to roughly $m_G \sim 1$ TeV.

The bad news is that for the reconstruction of the sgluon mass we could utilize the semileptonic sample. From Fig. [6] we can guess that the jet from the hadronic top decay should be comparable to the lepton spectrum, which does not guarantee that both of them even lie above $p_{T,j} > 50$ GeV, where the safe region would really only start at $p_{T,j} > 100$ GeV, as seen in Fig. [7]. Reconstructing the sgluon including a hadronic top decay is unlikely to succeed in a realistic QCD environment. As mentioned above, single sgluon production might rely on such a reconstruction for a side-bin analysis against the theoretically notorious single-top background, so QCD effects are unlikely to help with this problem.

V. CONCLUSIONS

The MRSSM, with a continuous $R$ symmetry, is an interesting alternative to the MSSM, which provides a novel solution to the tension between supersymmetry at the TeV scale and measurements from the flavor sector. A key feature of any $R$-symmetric model is the promotion of gauginos to Dirac fermions, which automatically implies the existence of an additional color-octet scalar with Standard-Model type $R$ charge, the sgluon. Sgluons necessarily have tree-level couplings to gluons, gluinos, and squarks, which further induce couplings to quarks and (single sgluon coupling to) gluons through loops.

Sgluons have a large color charge and can be copiously produced at the LHC. Their decays can include squarks, gluinos, gluons, and quarks. This last decay, through the large squark mixing which is the hallmark of the MRSSM, can violate flavor, leading to large branching ratios into a top and a light quark. Pair production with a subsequent decay to light-sign top quarks appears to be the most promising search channel at the LHC. We have computed the relevant LHC cross sections and branching ratios, with special focus on jet activity in this essentially background-free signature.

Properly simulated QCD effects turn out to be helpful with regard to the sgluon pair analysis: additional jet radiation in the signal will typically create more hard jets, adding to the two hard decay jets from the sgluon pair. While the QCD jets are too soft to be mistaken for decay jets they create a generally harder signal event, while jet radiation for the already softer top-pair background makes this
general feature even more prominent. In the light of this result, a search for sgluon pairs with a $t\bar{t}$ final state might be feasible.

For single sgluon production the QCD-induced background uncertainty is potentially dangerous. The obvious way to tell apart signal and background would be to look for a peak in the top–jet invariant mass, with a hadronically decaying top quark. However, QCD jet radiation can be expected to lead to a sizeable combinatorial background to the $W$ decay jets, which is worsened by the generically higher scale of the signal process.

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APPENDIX A: LOOP INDUCED COUPLINGS OF SGLUONS

In this appendix, we present the expressions for sgluon coupling to quarks and gluons induced at one loop. The results are presented in terms of the scalar integral functions of Passarino and Veltman [25], normalized such that the measure is $d^q q/(i\pi^2)$.

The sgluon coupling to quark $q^i$ and anti-quark $\bar{q}^j$ receives contributions from both right- and left-handed squarks. In the MRSSM, right- and left-handed squarks do not mix with each other, so even with large squark mixing, they form two distinct sectors. We are interested in two cases. The first has $m_i \gg m_j$, for which we can approximate,

$$f_{q,i}^R = \sqrt{2} \frac{m_{G}^2}{m_{G}^2 - m_i^2} \times \left\{ N_c \left[ B \left( m_i^2; m_{\bar{q}}, m_q \right) - B \left( m_G^2; m_{\bar{q}}, m_{\bar{q}} \right) + (m_{\bar{q}}^2 - m_q^2) C \left( m_i^2, 0, m_G^2; m_{\bar{q}}, m_q, m_{\bar{q}} \right) \right] \right.$$ \hspace{1cm} \begin{equation} \left. - \frac{1}{N_c} \left[ B \left( m_i^2; m_{\bar{q}}, m_q \right) - B \left( m_G^2; m_{\bar{q}}, m_q \right) + (m_{\bar{q}}^2 - m_q^2) C \left( m_i^2, 0, m_G^2; m_{\bar{q}}, m_q, m_{\bar{q}} \right) \right] \right\} \right) \end{equation} \hspace{1cm} \text{(A1)}$$

where the left-handed squarks contribute. And,

$$f_{q,i}^L = -\sqrt{2} \frac{m_{G}^2}{m_{G}^2 - m_i^2} \times \left\{ N_c \left[ B \left( m_i^2; m_{\bar{q}}, m_q \right) - B \left( m_G^2; m_{\bar{q}}, m_{\bar{q}} \right) + (m_{\bar{q}}^2 - m_q^2) + m_i^2 \right] C \left( m_i^2, 0, m_G^2; m_{\bar{q}}, m_q, m_{\bar{q}} \right) \right] \right.$$ \hspace{1cm} \begin{equation} \left. - \frac{1}{N_c} \left[ B \left( m_i^2; m_{\bar{q}}, m_q \right) - B \left( m_G^2; m_{\bar{q}}, m_q \right) + (m_{\bar{q}}^2 - m_q^2) \right] C \left( m_i^2, 0, m_G^2; m_{\bar{q}}, m_q, m_{\bar{q}} \right) \right\} \right) \end{equation} \hspace{1cm} \text{(A2)}$$

where the right-handed squarks are running in the loops. The chiral couplings are a consequence of eq. (2) and will be different for different admixtures of $G$ and $G^*$. The $f_{q,i}^{(L/R)}$ multiply $m_j$ and thus can
be neglected. The second case has $m_i = m_j$. The relevant quantities are $f_{q,i}^{(R/L)} - f_{q,j}^{(R/L)}$ given by,

$$f_{q,i}^{(R/L)} - f_{q,j}^{(R/L)} = \sqrt{2} \frac{m_G^2}{m_G^2 - 4m_i^2} \times \left\{ 2N_c \left[ B \left( m_i^2; m_{\tilde{g}}, m_{\tilde{q}} \right) - B \left( m_G^2; m_{\tilde{g}}, m_{\tilde{q}} \right) + \left( m_{\tilde{g}}^2 - m_{\tilde{q}}^2 + m_i^2 - \frac{m_G^2}{2} \right) C \left( m_i^2, m_i^2, m_G^2; m_{\tilde{g}}, m_{\tilde{q}}, m_{\tilde{q}} \right) \right] \right. $$

$$- \frac{2}{N_c} \left[ B \left( m_i^2; m_{\tilde{g}}, m_{\tilde{q}} \right) - B \left( m_G^2; m_{\tilde{q}}, m_{\tilde{q}} \right) - \left( m_{\tilde{q}}^2 - m_i^2 + m_i^2 \right) C \left( m_i^2, m_i^2, m_G^2; m_{\tilde{g}}, m_{\tilde{q}}, m_{\tilde{q}} \right) \right] \right\} \quad (A3)$$

where the right-handed squarks contribute to $f_L$ and vice-versa. The sgluon coupling to gluons for one squark flavor is

$$\lambda_g = 2\sqrt{2} \sum \tilde{q} \left[ m_{\tilde{q}L}^2 C \left( 0, 0, m_{\tilde{q}L}^2, m_{\tilde{q}L}, m_{\tilde{q}L} \right) - m_{\tilde{q}R}^2 C \left( 0, 0, m_{\tilde{q}R}^2, m_{\tilde{q}R}, m_{\tilde{q}R} \right) \right] \quad (A4)$$

summed over all of the $n_f$ squark flavors. From the naive computation, the color factor of this coupling includes the antisymmetric $f_{abc}$ as well as the symmetric $d_{abc}$. However, due to the symmetry structure of $\text{eq.}(A3)$ only $d_{abc}$ survives. For the same reason, the gluino loop contribution with its only color structure $f_{abc}$ cancels completely.

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[1] G. D. Kribs, E. Poppitz and N. Weiner, arXiv:0712.2030 [hep-ph].
[2] V. Ruhlmann-Kleider, Acta Phys. Polon. B 38, 705 (2007).
[3] L. J. Hall and L. Randall, Nucl. Phys. B 352, 289 (1991).
[4] see e.g.: W. Altmannshofer, A. J. Buras and D. Guadagnoli, JHEP 0711, 065 (2007); L. J. Hall, V. A. Kostelecky and S. Raby, Nucl. Phys. B 267, 415 (1986); J. S. Hagelin, S. Kelley and T. Tanaka, Nucl. Phys. B 415, 293 (1994); F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B 477, 321 (1996); M. Misiak, S. Pokorski and J. Rosiek, Adv. Ser. Direct. High Energy Phys. 15, 795 (1998); J. A. Casas and S. Dimopoulos, Phys. Lett. B 387, 107 (1996); Y. Nir, JHEP 0705, 102 (2007); for a recent collection of constraints see also S. Dittmaier, G. Hiller, T. Plehn and M. Spannowsky, Phys. Rev. D 77, 115001 (2008).
[5] A. E. Blechman and S. P. Ng, JHEP 0806, 043 (2008).
[6] P. J. Fox, A. E. Nelson and N. Weiner, JHEP 0208, 035 (2002); Z. Chacko, P. J. Fox and H. Murayama, Nucl. Phys. B 706, 53 (2005) [arXiv:hep-ph/0406142].
[7] S. D. L. Amigo, A. E. Blechman, P. J. Fox and E. Poppitz, arXiv:0809.1112 [hep-ph].
[8] B. A. Dobrescu, K. Kong and R. Mahbubani, JHEP 0707, 006 (2007).
[9] M. I. Gresham and M. B. Wise, Phys. Rev. D 76, 075003 (2007) [arXiv:0706.0909 [hep-ph]]; M. Gerbush, T. J. Khoo, D. J. Phalen, A. Pierce and D. Tucker-Smith, Phys. Rev. D 77, 095003 (2008); P. Fileviez Perez, R. Gavin, T. McElmury and F. Petriello, Phys. Rev. D 78, 115017 (2008) [arXiv:0809.2106 [hep-ph]].
[10] B. A. Dobrescu, K. Kong and R. Mahbubani, arXiv:0709.2378 [hep-ph].
[11] M. Bona et al. [UTfit Collaboration], JHEP 0803, 049 (2008); for a similar analysis see also: J. Charles et al. [CKMfitter Group], Eur. Phys. J. C 41, 1 (2005) [arXiv:hep-ph/0406184].
[12] J. Alwall et al., JHEP 0709, 028 (2007).
[13] using a modified version of: M. Spira, arXiv:hep-ph/9510347.
[14] D. Lillie, J. Shu and T. M. P. Tait, JHEP 0804, 087 (2008).
[15] S. Bar-Shalom, A. Rajaraman, D. Whiteson and F. Yu, arXiv:0803.3795 [hep-ph].
[16] B. W. Harris, E. Laenen, L. Phaf, Z. Sullivan and S. Weinzierl, Phys. Rev. D 66, 054024 (2002); J. Campbell, R. K. Ellis and F. Tramontano, Phys. Rev. D 70, 094012 (2004); Q. H. Cao, R. Schwienhorst and C. P. Yuan, Phys. Rev. D 71, 054023 (2005); S. Frizione, E. Laenen, P. Motylinski and B. R. Webber, JHEP 0603, 092 (2006).
[17] T. M. P. Tait and C. P. P. Yuan, Phys. Rev. D 63, 014018 (2001).
[18] M. L. Mangano, M. Moretti and R. Pittau, Nucl. Phys. B 632, 343 (2002); J. Alwall, P. Artoisenet, S. de Visscher, C. Duhr, R. Frederix, M. Herquet and O. Mattelaer, arXiv:0809.2410 [hep-ph].
[19] S. Catani, F. Krauss, R. Kuhn and B. R. Webber, JHEP 0111, 063 (2001); A. Schälicke and F. Krauss, JHEP 0507, 018 (2005).
[20] J. Alwall et al., Eur. Phys. J. C 53, 473 (2008).
[21] for a pedagogical user’s introduction see e.g. T. Plehn, arXiv:0810.2281 [hep-ph].
[22] T. Plehn, D. Rainwater and P. Skands, Phys. Lett. B 645, 217 (2007); J. Alwall, M. P. Le, M. Lisanti and J. G. Wacker, [arXiv:0809.3264] [hep-ph].

[23] see e.g. B. C. Allanach, C. G. Lester, M. A. Parker and B. R. Webber, JHEP 0009, 004 (2000); C. G. Lester, M. A. Parker and M. J. White, JHEP 0601, 080 (2006).

[24] B. K. Gjelsten, D. J. Miller and P. Osland, JHEP 0506, 015 (2005); D. J. Miller, P. Osland and A. R. Raklev, JHEP 0603, 034 (2006).

[25] G. Passarino and M. J. G. Veltman, Nucl. Phys. B 160, 151 (1979).

[26] S. Y. Choi, M. Drees, J. Kalinowski, J. M. Kim, E. Popenda and P. M. Zerwas, [arXiv:0812.3586] [hep-ph].