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PHENOMENOLOGY OF R-SYMMETRIC SUPERSYMMETRY*

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Salient features of the Minimal R-symmetric Supersymmetric Standard Model and its distinct phenomenological implications at colliders are discussed.

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1. Introduction

After disturbing BBC news "Researchers failed to find evidence of so-called 'supersymmetric' particles, which many physicists had hoped would plug holes in the current theory" dated August 27, 2011 [1] you may ask a question: why should I listen to another SUSY talk? It is true that so far no signal of SUSY has been seen at the LHC and the exclusion limits on SUSY parameters have been substantially advanced. However, even in the simplest and most constrained version of the Minimal Supersymmetric Standard Model (MSSM), the explored part of the parameter space is still rather small compared to what had been envisaged in the preparations for the LHC [2]. In less constrained models the exclusion limits are far less restrictive. Thus all the theoretical arguments in favor for TeV-scale supersymmetry are still valid [3].

2. *R*-symmetry and the MRSSM

Even before the recent LHC constraints the supersymmetric models were under pressure due to dangerous lepton- and baryon-number violating terms in the superpotential, as well as flavor- and CP-violating soft-susy breaking masses and couplings. An attractive possibility to remove

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all these phenomenologically embarrassing terms is provided by imposing an R-symmetry [4]. It is a continuous global U(1) symmetry under which the Grassmannian coordinates transform as $\theta \to e^{i\xi}\theta$. It implies that the component fields of the supersymmetric superfields differ by the R-charge. Assigning the R-charge as $R(\theta) = +1$ we get $R(d\theta) = -1$. Since the gauge vector superfields $\hat{G}$ are real, they must have $R(\hat{G}) = 0$, implying $R = 0$ for the gauge fields G^μ, and $R = 1$ for the gauginos $\tilde{G}^\alpha$. The kinetic terms are, therefore, automatically R-symmetric independently of the assigned R-charges of the chiral superfields. On the other hand, the terms in the superpotential must have $R = +2$ to provide the R-symmetric potential, while the soft-SUSY breaking terms must have $R = 0$. If we adopt the R-charges of the MSSM superfields as in Table I, the μ term and baryon- and lepton-number changing terms in the superpotential as well as soft-supersymmetry breaking Majorana gaugino masses and trilinear A-terms are forbidden. Thus the L- and R-squark and slepton mass mixing is absent, as well as dimension-five operators mediating proton decay, while Majorana neutrino masses can be generated. With such an assignment the R-charges of the SM particles are 0, while the R-charges of their superpartners are $R = \pm 1$ (resembling the situation in the MSSM with R-parity, where $R_\mathrm{p} = 1$ for the SM particles and $R_\mathrm{p} = -1$ for superpartners).

TABLE I

The R-charges of the MSSM and of the new gauge chiral $\hat{\Sigma}_K$ and R-Higgs superfields and their bosonic and fermionic component fields.

Field	Superfield	Boson	Fermion
Matter $\hat{L}$, $\hat{E}^c$, $\hat{Q}$, $\hat{D}^c$, $\hat{U}^c$	+1	+1	0
Gauge vector $\hat{G}_K$	0	0	+1
Higgs $\hat{H}_d$, $\hat{H}_d$	0	0	−1
Gauge chiral $\hat{\Sigma}_K$	0	0	−1
R-Higgs $\hat{R}_{d,u}$	+2	+2	+1

Since the gaugino Majorana mass terms, as well as the conventional Higgs-higgsino μ term, are forbidden, the superfield content of the model must be extended in order to give non-zero masses to the gauginos and higgsinos. Introducing new chiral superfields $\hat{\Sigma}_K = \{\sigma_K, \tilde{G}'_K\}$ in the adjoint representation of the gauge group (with $K = C, I, Y$ for SU(3), SU(2), U(1) respectively) in addition to the standard vector superfields $\hat{G}_K = \{\tilde{G}_K, G^\mu_K\}$, as well as two iso-doublet chiral superfields $\hat{R}_u$, $\hat{R}_d$ (R-Higgs) to complement the standard Higgses $\hat{H}_u$, $\hat{H}_d$ defines the so-called Minimal R-symmetric Supersymmetric Standard Model (MRSSM) [5,6]. An alternative formulation can be found in [7].

The assignment of R-charges, as in Table I, admits the trilinear Yukawa terms $y_d \hat{H}_d \cdot \hat{Q}\hat{D}^c$ *etc.* in the superpotential, while forbids the standard μ-term as well as L- and B-violating couplings. The presence of the new R-Higgs superfields $\hat{R}_{d,u}$ with $R = 2$, however, allows bilinear μ-type mass terms as well as trilinear terms for isospin and hypercharge interactions in the superpotential

$$\mathcal{W}_R = \mu_d \hat{H}_d \cdot \hat{R}_d + \mu_u \hat{H}_u \cdot \hat{R}_u , \tag{1}$$
$$\mathcal{W}'_R = \lambda_d^{I,Y} \hat{H}_d \cdot \hat{\Sigma}_{I,Y} \hat{R}_d + \lambda_u^{I,Y} \hat{H}_u \cdot \hat{\Sigma}_{I,Y} \hat{R}_u . \tag{2}$$

Both terms can be thought of as generated by the Giudice–Masiero mechanism, $\frac{1}{M} \int d^4\theta \, \hat{X}^\dagger \, \mathcal{W}^{(\prime)}$, when the hidden sector chiral spurion $\hat{X}$ with $R = 2$ develops the vacuum expectation value $\langle \hat{X} \rangle = \theta^2 F$ [8].

In a similar way the soft-supersymmetry breaking parameters can also be generated by the R-symmetric interaction. The B_μ term can be generated as: $\frac{1}{M^2} \int d^4\theta \, \langle \hat{X}^\dagger \hat{X} \rangle \; \hat{H}_u \cdot \hat{H}_d \; \rightarrow B_\mu H_u \cdot H_d$. No bilinear coupling of the R-Higgs is allowed by the R-symmetry, thus breaking the exchange symmetry between the H and R-Higgs fields. The soft R-Higgs masses are generated similarly, *e.g.*: $\frac{1}{M^2} \int d^4\theta \, \langle \hat{X}^\dagger \hat{X} \rangle \; \hat{R}_d^\dagger \hat{R}_d \; \rightarrow -m^2_{R_d}(|R_d^+|^2 + |R_d^0|^2)$. The Dirac gaugino masses can be generated by the interaction of the form $\frac{1}{M} \int d^2\theta \, \langle \hat{W}'^\alpha \rangle \; \mathrm{Tr}\, \hat{G}_\alpha \hat{\Sigma} \; \rightarrow -M^D \tilde{G}\tilde{G}'$ between the gauge vector and adjoint chiral superfields with a hidden sector U(1) gauge superfield $\hat{W}'^\alpha$ with $R = 1$ which develops a vacuum D-term $\langle \hat{W}'^\alpha \rangle = D\, \theta^\alpha$ ($\hat{G}_\alpha$ are the gauge superfield-strengths with $R = 1$). Likewise, the bilinear couplings of $\hat{W}'^\alpha$ and $\hat{X}$ fields with the adjoint chiral superfields: $\frac{1}{M^2} \int d^2\theta \, \langle \hat{W}'^\alpha \hat{W}'_\alpha \rangle \, \mathrm{Tr}\, \hat{\Sigma}^2$, $\frac{1}{M^2} \int d^4\theta \, \langle \hat{X}^\dagger \hat{X} \rangle \, \mathrm{Tr}\, \hat{\Sigma}^2$, $\frac{1}{M^2} \int d^4\theta \, \langle \hat{X}^\dagger \hat{X} \rangle \, \mathrm{Tr}\, \hat{\Sigma}^\dagger \hat{\Sigma}$ generate soft masses of the adjoint scalars (M denotes a generic scale parameter which can be different in the above interactions).

3. Phenomenology of the MRSSM

The transition from Majorana gauginos to Dirac gauginos as well as presence of new $R_{u,d}$ and adjoint scalar fields σ has far reaching consequences on supersymmetric particle production at the LHC and e^+e^- colliders [9, 10, 11], cold dark matter expectations [12], and flavour- and CP-changing processes [5, 13]. The production and decays of colored scalars (sgluons) at the LHC is discussed in the talk by Kotlarski [14]; see also [15]. Here we will discuss how the nature of gauginos can be tested and what are the expectations for producing the R-Higgs bosons at colliders.

3.1. Nature of gauginos

With the Majorana mass terms forbidden, the two Majorana gluinos $\tilde{G}_C \equiv \tilde{g}$ and $\tilde{G}'_C \equiv \tilde{g}'$ are coupled by the supersymmetry-breaking Dirac mass. Therefore, they can be combined into a single Dirac fermion field $\tilde{g}_{\rm D} = \tilde{g}_{\rm R} + \tilde{g}'_{\rm L}$ with R-charge $+1$ (where $\tilde{g}_{\rm L,R} = \frac{1}{2}(1 \mp \gamma_5)\tilde{g}$, $\tilde{g}$ are matrix-valued $\tilde{g} = \frac{1}{\sqrt{2}}\lambda^a \tilde{g}^a$, and similarly for $\tilde{g}'$). Note that $g_{\rm D}$ is not self-conjugate: $\tilde{g}^c_{\rm D} \neq \pm\tilde{g}^T_{\rm D}$, and the anti-gluino $\tilde{g}^c_{\rm D}$ carries $R = -1$ charge, while $\tilde{g}^c_{\rm L} = -\tilde{g}^T_{\rm R}$ and $\tilde{g}'^c_{\rm L} = -\tilde{g}'^T_{\rm R}$ as required for Majorana fields [10]. Similarly, in the neutral electroweak sector the original MSSM gaugino $\tilde{G}$ and higgsino $\tilde{H}$ are coupled to the new gaugino $\tilde{G}'$ and higgsino $\tilde{R}$ fields by the Dirac mass terms as well as terms coming from bilinear and trilinear interactions, Eqs. (1), (2), giving rise to four mass-eigenstate Dirac neutralinos $\tilde{\chi}^0_{\rm D1,\dots,D4}$ [6]. Since all supersymmetric fermions are Dirac-type, it is convenient to introduce a conserved quantum number D as

$$
\begin{aligned}
D[\tilde{q}_{\rm L}] &= D[\tilde{g}^c_{\rm D}] = D[\tilde{l}^-_{\rm L}] = D[\tilde{\chi}^{c0}_{\rm D}] = D[\tilde{\chi}^+_{D1}] = D[\tilde{\nu}_{\rm L}] = -1\,, \\
D[\tilde{q}_{\rm R}] &= D[\tilde{g}_{\rm D}] = D[\tilde{l}^-_{\rm R}] = D[\tilde{\chi}^0_{\rm D}] = D[\tilde{\chi}^+_{D2}] = +1
\end{aligned}
\tag{3}
$$

with their antiparticles carrying the opposite D-charge; the D-charge for the SM particles vanishes.

The conserved D-charge then easily allows to identify the processes that are forbidden in the Dirac case in contrast to the Majorana. For example, the processes $qq' \to \tilde{q}_{\rm L}\tilde{q}'_{\rm L}$, $qq' \to \tilde{q}_{\rm R}\tilde{q}'_{\rm R}$ allowed in the Majorana gluino case are forbidden in the Dirac theory, while $qq' \to \tilde{q}_{\rm L}\tilde{q}'_{\rm R}$ are allowed in both (and equal). Likewise, $q\bar{q}' \to \tilde{q}_{\rm L}\tilde{q}'^*_{\rm R}$ production allowed for the Majorana gluinos is forbidden in the Dirac, while $q\bar{q}' \to \tilde{q}_{\rm L}\tilde{q}'^*_{\rm L}$ and $q\bar{q} \to \tilde{q}_{\rm R}\tilde{q}'^*_{\rm R}$ are allowed in both (and equal). In a similar manner one easily finds which processes are allowed in the quark–gluon and gluon–gluon processed. At the LHC the $\tilde{q}_{\rm L}$ and $\tilde{q}_{\rm R}$ are not easily distinguishable. Nevertheless, owing to the valence quark distributions in the proton beams and different decay modes of $\tilde{q}_{\rm L}$ and $\tilde{q}_{\rm R}$, the Dirac–Majorana nature of gluinos could be resolved, *e.g.* by counting numbers of same-sign and opposite-sign lepton pairs, as discussed at length in [9, 10].

The nature of neutralinos could also be checked at the LHC if long decay chains of squarks with intermediate neutralinos and sleptons $\tilde{q}_{\rm L} \to q\tilde{\chi}^0_i \to l_{\rm n}\tilde{l} \to ql_{\rm n}l_{\rm f}\tilde{\chi}^0_1$ are identified ($l_{\rm n}$ denotes the *near* lepton from neutralino decay, and $l_{\rm f}$ is the *far* one from slepton decay). Neutralinos produced in $\tilde{q}_{\rm L}$ decays are preferentially left-handed. If only singlet slepton $\tilde{l}_{\rm R}$ is accessible in the $\tilde{\chi}^0_i$ decay, then the near lepton is either left-handed l^+ or right-handed l^-. Angular momentum conservation then implies that l^+ (l^-) is preferentially flying in parallel (opposite) to the neutralino direction, *i.e.* in the rest frame of $\tilde{q}_{\rm L}$ a near l^+ will tend to be harder than l^-. The same arguments show that a far l^+ will also tend to be harder. Such correlations find their imprint in the ql mass distributions [16]. In the case of Majorana neutralinos the $\tilde{\chi}_i^0$ can decay to both $\tilde{l}_{\rm R}^{\pm}$, $\tilde{q}_{\rm L} \to q\tilde{\chi}_i^0 \to l_{\rm n}^{\mp}\tilde{l}_{\rm R}^{\pm} \to ql_{\rm n}^{\mp}l_{\rm f}^{\pm}\tilde{\chi}_1^0$, thus the ql^+ (ql^-) mass distribution gets contributions from near and far l^+ (l^-), each coming with endpoints determined by the masses of sparticles in the corresponding decay chain. The distributions are the same for squark and anti-squark decays. In the case of Dirac, only the $\tilde{q}_{\rm L} \to q\tilde{\chi}_{{\rm D}i}^{c0} \to ql_{\rm n}^-\tilde{l}_{\rm R}^+ \to ql_{\rm n}^-l_{\rm f}^+\tilde{\chi}_{\rm D1}^{c0}$ is allowed, while for the anti-squark $\tilde{q}_{\rm L}^* \to \bar{q}\tilde{\chi}_{{\rm D}i}^0 \to \bar{q}l_{\rm n}^+\tilde{l}_{\rm R}^- \to \bar{q}l_{\rm n}^+l_{\rm f}^-\tilde{\chi}_{\rm D1}^0$. Therefore, the ql^+ mass distribution will be different for both cases, and distinctly different from the Majorana case, as illustrated in the left panel of Fig. 1 (taken from [10] calculated for sparticle masses according to the SPS1a$'$ scenario [17]).

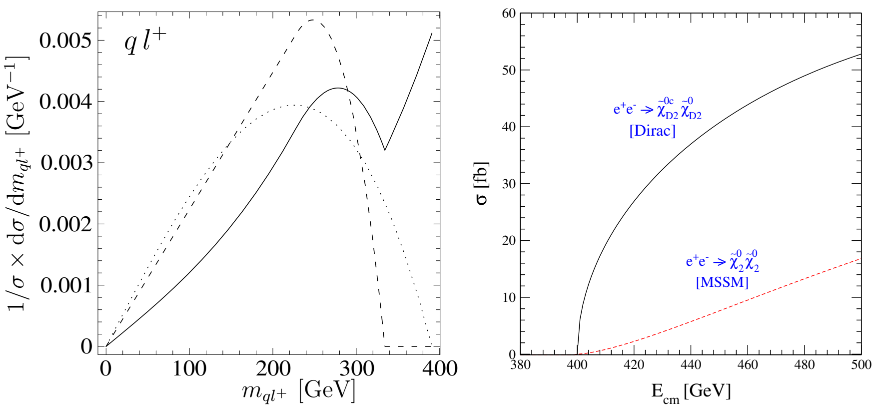

Fig. 1. Left: ql^+ invariant mass distributions for squark and anti-squark decay chains in the Majorana case (solid), and squark (dashes) and anti-squark (dots) for Dirac. Right: near threshold behavior of the diagonal neutralino pair production cross-section in e^+e^- collisions. See text for details.

An $e^{\pm}e^-$ collider with polarized beams is ideally poised for testing the nature of neutralinos. In parallel to the quark/squark case, the processes $e^-e^- \to \tilde{e}^-\tilde{e}^-$ with equal chiralities $(\tilde{e}_{\rm L}^-\tilde{e}_{\rm L}^-$ or $\tilde{e}_{\rm R}^-\tilde{e}_{\rm R}^-)$ and $e^+e^- \to \tilde{e}_{\rm L}^+\tilde{e}_{\rm R}^-$ are forbidden by the conserved D, while in general non-zero in the Majorana case. Another way to test the nature of neutralinos is provided by the near-threshold behavior of the diagonal neutralino pair production and their angular distribution. For the Dirac particles the cross-section for $e^+e^- \to \tilde{\chi}_{{\rm D}i}^c\tilde{\chi}_{{\rm D}i}^0$ shows a typical s-wave excitation (and non F-B symmetric), while for Majorana the diagonal pair is excited in the p-wave, as shown in the right panel of Fig. 1 [11].

3.2. R-Higgs bosons

The non-colored scalar potential in general is very complex due to the mixing among the H, R and $\sigma_{I,Y}$ states. Since the mass parameters of the SU(2) sector must be large, as demanded by the ρ parameter, we assume that the mass parameters in the $\sigma_{I,Y}$ sectors are large, of TeV order[1]. Leaving out the $\sigma_{I,Y}$ fields, the neutral part of the potential simplifies to

$$\begin{aligned}\mathcal{V}^0_{[H,R]} \;=\;& \left(m^2_{H_d}+\mu^2_d\right)\left|H^0_d\right|^2+\left(m^2_{H_u}+\mu^2_u\right)\left|H^0_u\right|^2-\left(B_\mu H^0_d H^0_u+\text{h.c.}\right)\\ &+\left(m^2_{R_d}+\mu^2_d\right)\left|R^0_d\right|^2+\left(m^2_{R_u}+\mu^2_u\right)\left|R^0_u\right|^2\\ &+\left|\lambda^I_d H^0_d R^0_d+\lambda^I_u H^0_u R^0_u\right|^2+\left|\lambda^Y_d H^0_d R^0_d-\lambda^Y_u H^0_u R^0_u\right|^2\\ &+\tfrac{1}{8}\left(g^2+g'^2\right)\left(|H^0_d|^2-|H^0_u|^2-|R^0_d|^2+|R^0_u|^2\right)^2 . \end{aligned}\tag{4}$$

The absence of the mixed $R^0_d R^0_u$ term, as required by R-symmetry, implies that the R-Higgs fields, $R_{d,u}$, do not develop non-zero vacuum expectation values and, as a result, do not mix with the H fields; this conclusion is still valid even if the $\sigma_{I,Y}$ fields are not neglected.

Since the same-sign charged $R_{d,u}$-Higgs bosons carry opposite R-charges, they do not mix, while the mass-eigenstates of the neutral ones are obtained by a standard diagonalization procedure. Thus, the R-symmetric model possesses four neutral and four charged R-Higgs states in addition to the standard three neutral and two charged MSSM H-Higgs states. For $(m^2_{R_{d,u}}+\mu^2_{d,u})^{1/2}\equiv m'_{\rm R}\geq v$ the neutral and charged masses are roughly equal to $m'_{\rm R}$ *modulo* terms of order $g^2v^2/m'_{\rm R}$, in analogy to the heavy Higgs bosons of the MSSM.

The conserved R-charge restricts the R-Higgs boson trilinear couplings to pairs of sfermions, $R\tilde{\ell}\tilde{\ell}$, $R\tilde{q}\tilde{q}$ and chargino/neutralino combinations, $R\tilde{\chi}\tilde{\chi}$; the couplings to pairs of SM particles and Higgs bosons, Rff, RVV, RHH vanish (even at loop order). The R-symmetry admits also RR^*H and RR^*V couplings. Thus in pp and e^+e^- collisions the R-Higgs bosons can be produced in pairs. Fig. 2 shows the expected size of the cross-sections for the production of neutral/charged R-Higgs pairs at the LHC and e^+e^- colliders; cross-sections for the diagonal neutral R-Higgs boson pairs, $R^0_1 R^{0^*}_1$ and $R^0_2 R^{0^*}_2$, vanish for the common R-Higgs mass parameter $m'_{\rm R}=(m^2_{R_{d,u}}+\mu^2_{d,u})^{1/2}$ (the $\lambda^{I,Y}$ couplings are taken as predicted by $N=2$ susy). Additional sources of R Higgs bosons, though in general at reduced levels, are provided by the fusion channel $pp\to\gamma\gamma\to R^+R^-$, or from heavy MSSM Higgs decays $H\to RR^*$.

[1] Sgluons σ_C can nevertheless be lighter and produced copiously at the LHC, see [14].

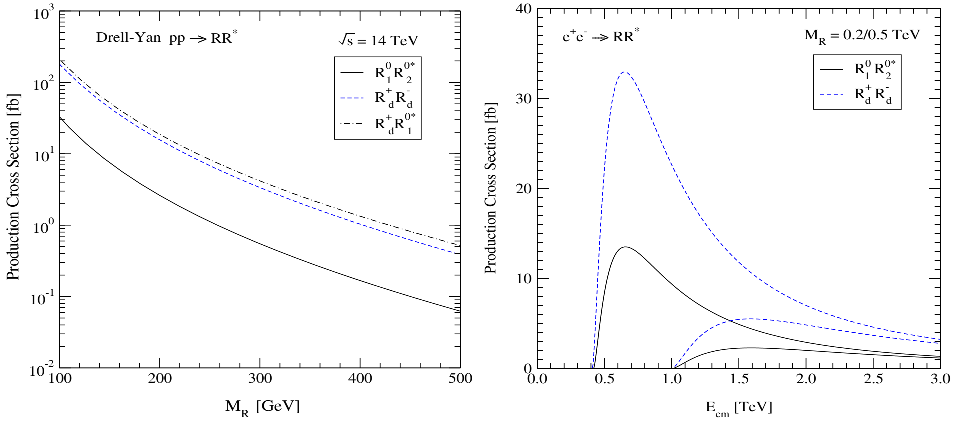

Fig. 2. Left: Drell–Yan production cross-sections of R-Higgs boson pairs at the LHC ($\sqrt{s} = 14$ TeV) as a function of the averaged mass of the produced particles $M_{\rm R}$; Right: Production of t R-Higgs boson pairs at e^+e^- colliders [ILC/CLIC] for $M_{\rm R} = 0.2$ TeV and 0.5 TeV.

Although the details of the experimental signature depend on the specific scenario, the pair-production of R-Higgs bosons determines the main characteristics. Since we assume $M_{\rm R} > 2M_{\tilde{\chi}^0_{D1}}$, the R-Higgs bosons are unstable and decay to pairs of sfermions/charginos (which further decay) and pairs Dirac neutralinos. Taking the SPS1a$'$ scenario as an illustration, the τ's are the dominating visible cascade components, *cf.* Ref. [18]: $R^0 \to \tilde{\chi}^0_{\rm D1}\tilde{\chi}^0_{\rm D2}$ followed by $\tilde{\chi}^0_{\rm D2} \to \tau\tilde{\tau}$ followed by $\tilde{\tau} \to \tau\tilde{\chi}^0_{\rm D1}$. In such a case one would expect

$$R^0 R^{0^*} \to \tau^+\tau^-\tau^+\tau^- + \tilde{\chi}^0_{\rm D1}\tilde{\chi}^0_{\rm D1}\tilde{\chi}^{0c}_{\rm D1}\tilde{\chi}^{0c}_{\rm D1}\,. \tag{5}$$

The multi-fold τ-multiplicity in association with high values of missing energy/transverse momentum offers promising signatures for detecting RR events.

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