R Parity from the Heterotic String

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In T-duality invariant effective supergravity with gaugino condensation as the mechanism for supersymmetry breaking, there is a residual discrete symmetry that could play the role of R parity in supersymmetric extensions of the standard model.

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One of the challenges of string theory is to provide a mechanism for forbidding operators that violate lepton number and baryon number in the low energy effective theory. In the minimal supersymmetric extension (MSSM) of the standard model, this is achieved by imposing a discrete symmetry, called R parity, such that the unwanted operators are forbidden, while those that give masses to quarks and leptons are allowed, as is the Higgs boson mass term (“μ term”) that is needed to produce the correct electroweak symmetry breaking pattern.

In the context of the weakly interacting heterotic string, a favored mechanism for breaking supersymmetry uses condensation that occurs in a hidden sector when some gauge group becomes strongly coupled. These theories include gauge invariant scalar fields, notably T moduli whose vacuum values (VEVs) determine the size of six compact dimensions, and the dilaton whose VEV determines the gauge coupling constant at the string scale. They are invariant [1] under a discrete symmetry called T duality, which is broken when the T moduli are stabilized at their VEVs. A variety of effective supergravity Lagrangians with these properties have been constructed; in a class [2,3] of these where (nearly) vanishing vacuum energy is imposed, T duality assures that the T moduli are stabilized at self-dual points, defined below, with vanishing VEVs for their auxiliary fields. Supersymmetry breaking is dilaton mediated; this avoids a potentially dangerous source of flavor changing neutral currents (FCNC), since the dilaton couplings to matter are universal, while the T-moduli couplings are not. Another consequence of this result is that there is a residual discrete symmetry that might play the role of R parity in the MSSM. This Letter explores the discrete symmetries of a particular class of string compactifications. [This analysis is not expected to be applicable in many models (for example, “racetrack” models) where the T moduli are stabilized away from self-dual points. Aside from potential difficulties with FCNC, these models need to invoke a second source of supersymmetry breaking to cancel the cosmological constant. The effects of quadratically divergent loop corrections [4] on FCNC and the cosmological constant will be examined elsewhere [5].]

In the class of models considered here, transformations on the T moduli take the form

$$\begin{gathered} T^I \to \frac{a^I T^I - ib^I}{ic^I T^I + d^I}, \qquad \Phi^A \to e^{-\sum_I q_I^A F^I}\Phi^A, \\ F^I = \ln(ic^I T^I + d^I), \qquad a^I d^I - b^I c^I = 1, \\ a^I, b^I, c^I, d^I \in \mathbf{Z} \quad \forall\ I = 1, 2, 3, \end{gathered} \tag{1}$$

and the Kähler potential K and superpotential W transform as $K \to K + F + \bar{F}$, $W \to e^{-F}W$, $F = \sum_I F^I$. (We neglect mixing [6] among twisted sector fields of the same modular weights q_I^A with mixing parameters that depend on the integers a^I, b^I, c^I, d^I.) The self-dual vacua $T_{\rm sd}$, namely, $\langle t^I\rangle = 1$ or $e^{i\pi/6}$, are invariant under (1) with

$$b^I = -c^I = \pm 1, \qquad a^I = d^I = 0 \quad \text{or} \quad \begin{cases} a^I = b^I, & d^I = 0, \\ d^I = c^I, & a^I = 0, \end{cases} \qquad F^I = ni\frac{\pi}{2} \quad \text{or} \quad ni\frac{\pi}{3}. \tag{2}$$

For three moduli there is a symmetry under $G_{\rm sd} = Z_2^m \otimes Z_3^{m'}$, $m + m' = 3$. The gaugino and matter condensates that get VEVs break this further to a subgroup G_R with $i\,{\rm Im}F = F = 2ni\pi$, under which $\lambda_L \to e^{-(i/2)\,{\rm Im}F}\lambda_L = \pm\lambda_L$; we would identify the case with a minus sign with R parity. This subgroup also leaves invariant the soft supersymmetry-breaking terms in the observable sector, if no other field gets a VEV that breaks it. For example, if the μ term comes from a superpotential term $H_u H_d \Phi$, with the VEV $\langle \phi = \Phi| \rangle \neq 0$ generated at the TeV scale, the symmetry could be broken further to a subgroup $R \in G_R$ such that $R\Phi = \Phi$. On the other hand, if the μ term comes from a Kähler potential term generated by invariant VEVs above the scale where the moduli are fixed, there would be no further breaking down to the electroweak scale.

Superpotential terms of the form $W = \prod_A \Phi^A \times \prod_I \eta(iT^I)^{2(\sum_A q_I^A - 1)}$ would be covariant under (1), given the transformation property of the Dedekind η function,

$$\begin{gathered} \eta(iT^I) \to e^{i\delta_I} e^{(1/2)F(T^I)}\eta(iT^I), \qquad F(T^I) = F^I, \\ \delta_I = \delta_I(a^I, b^I, c^I, d^I), \end{gathered} \tag{3}$$

if the moduli independent phases [7] satisfied $\delta_I = 2n^I i\pi$. It is easy to see that this is not the case for the transformations that leave fixed the self-dual points $T_{\rm sd}$: $\eta(iT_{\rm sd}) \to \eta(iT_{\rm sd}) \neq e^{(1/2)F(T_{\rm sd})}\eta(iT_{\rm sd})$. It follows from T duality that this phase can be reabsorbed [7] into the transformation properties of the twisted sector fields. Consider, for example, a Z_3 orbifold with untwisted sector fields U^{AI}, and twisted sector fields T^A and Y^{AI} with modular weights

$$(q_I^{AJ})_U = \delta_I^J, \qquad q_I^A = \left(\frac{2}{3}, \frac{2}{3}, \frac{2}{3}\right), \qquad (q_I^{AJ})_Y = \left(\frac{2}{3}, \frac{2}{3}, \frac{2}{3}\right) + \delta_I^J. \tag{4}$$

Since allowed superpotential couplings are of the form [8] $T^{3p}U^q$, invariance under (1) and (3) with $F = 2ni\pi$ requires that $\prod \Phi^A$ in W gets, in addition to the phases implicit in (1), an overall phase factor $\delta_{\rm sd}$ that satisfies

$$i\delta_{\rm sd} = \sum_I (2p + n^I)F(T_{\rm sd}^I) = \sum_I n^I F(T_{\rm sd}^I) + 4ik\pi, \qquad k \in \mathbf{Z}, \tag{5}$$

where n^I is the number of U^{AI} and T^{AI} factors in $\prod \Phi^A$. More generally, T duality implies that the allowed terms in the superpotential must be such that there is a choice of phases δ^A that makes it covariant if the transformation of Φ^A in (1) is modified to read

$$\Phi^A \to e^{i\delta^A - \sum_I q_I^A F^I}\Phi^A. \tag{6}$$

For example, in [7] trilinear terms $W \sim U_1U_2U_3$, $(T)^3$, where $T = T^A$, were considered; all such terms would be covariant provided $\delta^U = 0$, $\delta^T = -\frac{2}{3}\delta = -\frac{2}{3}\sum_I \delta_I$. If we further impose $\delta^{Y^J} = -\frac{2}{3}\delta - 4\delta_I$, then the monomials

$$U_1U_2U_3, \qquad T^3\Pi^p, \qquad U_IY^IT^2\Pi^p, \qquad U_IY^IU_JY^JT\Pi^p, \qquad U_IY^IU_JY^JU_KY^K\Pi^p, \tag{7}$$

where $\Pi = Y^1Y^2Y^3$, can be used to construct covariant superpotential terms W_α, by multiplying them by powers of $\eta_I = \eta(iT^I)$ such that the overall monomials W_α are modular covariant with modular weights $q_I^\alpha = 1$. Further operators can be constructed by multiplying these by invariant operators that can also appear in the Kähler potential; for example, $\Pi\eta^6$, $\eta = \prod_I \eta_I$, is invariant. In addition, invariant operators of the form $\eta^{2m}\prod_{i=1}^m W_{\alpha_i}$, $2m\delta = 2\pi n$, can be constructed since δ/π is a rational number. [The group (1) of duality transformations on T^I is generated [9] by $T^I \to 1/T^I$ with $\delta(0, 1, -1, 0) = \pi/4$, and $T^I \to T^I - i$ with $\delta(1, 1, 0, 1) = \pi/12$.] These couplings are consistent with the selection rules [8]. They are further restricted by additional selection rules and gauge invariance. For the subgroup defined by (2) and $F = 2ni\pi$, $i\delta = -\frac{1}{2}F = -in\pi$, the superpotential is invariant, as are the monomials in (7), so any product of them in could appear in the effective superpotential or Kähler potential, e.g., through quantum corrections and/or integrating out massive fields, in the effective theory below the scale where the T moduli are fixed and supersymmetry is broken, with possibly additional VEVs that are invariant under G_R generated at that scale.

Superpotential terms of dimension three will be generated from higher order terms when some fields acquire VEVs. In models with an anomalous $U(1)_X$, there is a Green-Schwarz counterterm in the form of a D term [10] that leads to the breaking of a number m of $U(1)$ gauge factors when $n \geq m$ fields Φ^A acquire VEVs. T duality remains unbroken [11], but the modular weights are modified by going to unitary gauge in a way that keeps modular invariance manifest. For example, in minimal models with $n = m$,

$$\Phi^M \to \Phi'^M, \qquad q_I^M \to q_I'^M = q_I^M - \sum_{Aa} q_a^M Q_A^a q_I^A, \qquad \sum_A Q_A^a q_b^A = \delta_b^a, \qquad \sum_a Q_A^a q_a^B = \delta_A^B. \tag{8}$$

This has the effect of making the eaten chiral supermultiplets modular invariant in the super-Higgs mechanism. Then, for a term in the superpotential with some $\langle\Phi^A\rangle \neq 0$,

$$\begin{aligned} W &= \prod_M \Phi^M \prod_B \langle\Phi^A\rangle \prod_I \eta(iT^I)^{2(\sum_M q_I^M + \sum_B q_I^A - 1)}, \\ &= \prod_M \Phi'^M \prod_B \langle\Phi^A\rangle \prod_I \eta(iT^I)^{2(\sum_M q_I'^M - 1)}, \end{aligned} \tag{9}$$

because W is also $U(1)_a$ invariant: $\sum_M q_a^M + \sum_A q_a^A = 0$. In order to make T duality fully manifest below the $U(1)$-breaking scale, we have to redefine the transformation (6) by including a global $U(1)_a$ transformation such that Φ^A is fully invariant, and

$$\Phi'^M \to e^{i\delta'^M - \sum_I q_I'^M F^I}\Phi^M, \qquad \delta'^M = \delta^M - \sum_{Aa} q_a^M Q_A^a \delta^A. \tag{10}$$

A priori, we expect that $\langle\Phi^A\rangle \sim 0.1$, so that couplings arising from high dimension operators in the superpotential are suppressed. (The factors multiplying these terms can in fact be rather large [12].) We would like to have one large coupling (Q_3, T^c, H^u) which should correspond to one of the dimension three operators in (7). Most models studied [13,14] have quark doublets in the untwisted sector. In this case, we should take T^c and H^u in the untwisted sector as well, and require $q_a^{Q_3} + q_a^{T^c} + q^{H^u} = 0$. [This requirement in satisfied in the Font-Ibáñez-Quevedo-Sierra (FIQS) model [14].] That is, if we identify the Q_I generation index with the moduli index, we can have, e.g., $T^c = T_2^c$, $H^u = H_1^u$. Then to suppress the $Q_2C^cH^u$ and $Q_2U^cH^u$ couplings we require $C^c, U^c \notin U_3$, so one of these must be in the untwisted sector T. Since these generally have different $U(1)$ charges from the untwisted sector fields, to avoid a possible D-term induced flavor dependence of the squark masses in the first two generations, we also take both U^c and C^c in T.

As an example (that turns out not to produce the desired R parity) consider the FIQS model [14], with the ϕ^A vacuum studied in [3]. Then $D^c, S^c, B^c, H^u \in T$. To generate all the known Yukawa's (QT^2, QH^uT) it follows from (7) that at least three Y^I with different indices I have to have VEVs. If all the VEVs are generated by D-term breaking, we have to choose the threefold version [3] of the "minimal" FIQS model with $\langle Y_1^{1,2,3}\rangle \neq 0$. Defining $\zeta_A^M = \sum Q_A^a q_a^M$, $\zeta^M = \sum_A \zeta_A^M$, we have $\zeta_{Y_1}^M = \frac{2}{3}\zeta^M - \sqrt{\frac{3}{2}} q_X^M$. Then after the redefinitions (8) and (10) with $i\delta_I = -\frac{1}{2}F^I$, we have $T'^M \to e^{[1/3(\zeta^M-1)+\sqrt{3/2}q_X^M]F}T'^M$. For all the possible MSSM candidates the T^M have $\sqrt{\frac{3}{2}}q_X^M = \frac{1}{3}n^M$ and we just get $T'^M \to e^{1/3(\zeta^M-1+n^M)F}T'^M$. Then, using the constraint $F = 2ni\pi$, trilinear terms T^3 with fixed n^M satisfy $\prod_{i=1}^3 T'^{M_i} \to e^{1/3\sum_{j=1}^3 \zeta^{M_j}F}\prod_i T'^{M_i}$. Taking the T^{M_i} to be the FIQS supermultiplets $U^c = u_2 \in T$ and any two of $D^c, S^c, B^c = d_{1,2} \in T$, , we have $n^{M_i} = -1$, $\sum_j \zeta^{M_j} = 0$, so we cannot forbid baryon number violating couplings. Nevertheless, we can ask if we get any interesting restrictions. It turns out that for MSSM gauge invariant trilinear couplings we get $\sum_j \zeta^{M_j\neq \ell_5} = n$, $\sum_j \zeta^{M_j\neq\ell_5} + \zeta^{\ell_5} = \frac{n}{2}$, and after imposing $F = ni\pi$, aside from couplings involving ℓ_5, everything drops out except the original T-duality transformation on the untwisted fields; if for an operator O, $O \to \eta(O)O$, we have

$$\eta(G_iG_j\ell_{k\neq 5}) = 1, \qquad \eta(G_iG_j\ell_5) = e^{(n/3)i\pi}, \qquad \eta(Q_Id_iG_j) = e^{-F^I}, \qquad \eta(Q_Iu_J^1\tilde{G}_K^1) = e^{-F^I-F^J-F^K},$$
$$\eta(Q_Iu_2\tilde{G}_{k\neq 1}) = e^{-F^I}, \qquad \eta(Q_Iu_J^1\tilde{G}_{k\neq 1}) = \eta(Q_Iu_2\tilde{G}_J^1) = e^{-F^I-F^J}. \tag{11}$$

Thus L^2E^c is allowed unless $E^c = \ell_5$, in which case LH_dE^c is also forbidden, unless the symmetry is broken to $n = 6p$, in which case both are allowed. In order to have at least one $Q_IH_dD^c$-type coupling for each Q_I, we need $F^I = 2in\pi\forall I$. Then all couplings involving the Q_I are allowed, including Q_ILD^c, etc. We can also look at candidate μ-term couplings $H_uH_d = \tilde{G}_iG_j$; these have $\zeta^{\tilde{G}_i} + \zeta^{G_j} = m$, and

$$\eta(G_i\tilde{G}_I^1) = e^{(m/3)i\pi+F^I}, \qquad \eta(G_i\tilde{G}_{j\neq 1}) = e^{(2m/3)i\pi}, \tag{12}$$

so G_R has to be broken to a smaller subgroup when the μ term is generated.

Apart from the fact that the FIQS model does not give the correct constraints, it is still interesting to see if one can get any unbroken symmetry after the Higgs particles acquire VEVs. In this model, the individual η's are of the form $e^{2ni\pi(m/33)}$, except for ℓ_5 where $\frac{m}{33} \to \frac{2m+1}{66}$, so the individual VEVs of $H_{u,d}$ break the symmetry down to a subgroup with $F = 33pi\pi$. If this is the only symmetry left, the only constraint on the couplings in (11) is to forbid $G_iG_j\ell_5$. However, we can once again redefine the transformations such that one Higgs boson is invariant and the other has the phase factor in (12), and therefore both Higgs bosons are invariant under the subgroup left unbroken by the μ term. Here we use the fact that the couplings are invariant under electroweak hypercharge Y, and redefine the transformation properties by $\eta_M \to \eta_M\eta_{H_u}^{-2Y^M}$. Then H_u with $Y^{H_u} = \frac{1}{2}$ is invariant and the couplings that were allowed or forbidden under the group left unbroken by the μ term (12) remain so.

Now we turn to a more general analysis, assuming the same assignments as before for the MSSM fields, but with different $U(1)$ charges. Then under G_R

$$U_J'^M \to e^{(1/3)\zeta^M F - \sum_I \zeta_I^{MJ}F^I - F^J} U_J'^M = \eta_{MJ}U_J'^M,$$
$$T'^M \to e^{(1/3)(\zeta^M-1)F - \sum_I \zeta_I^M F^I} T'^M = \eta_M T'^M, \tag{13}$$

where $\zeta_I^M = \sum_A \zeta_{Y_I^A}^M$, $\zeta_A^Q + \zeta_A^{T^c} + \zeta_A^{H^u} = 0$, $\eta_{Q_3}\eta_{T^c}\eta_{H^u} = 1$. We also require $\eta_{H_u}\eta_{H_d} = 1$. If $\tilde{D}^c = D^c, S^c, B^c$ all have the same $U(1)$ charges, then $\eta_{D^c} = \eta_{S^c} = \eta_{B^c}$. Then in order to have at least one coupling $Q_I\tilde{D}^cH_d$ for each value of I, we require $\eta_{Q_I} = \eta_Q$ independent of I, and requiring at least one coupling $Q_IC^cH_u$ implies $\eta_{U^c} = \eta_{C^c} = \eta_{T^c}$ if $\tilde{U}^c = U^c, C^c$ have the same $U(1)$ charges. Then all couplings of these two types are allowed. Similarly, if we assume the lepton doublets L and singlets E^c have (two sets of) degenerate $U(1)$ charges $q_a^L, q_a^{E^c}$, they also have degenerate R parities: η_L, η_{E^c}. To forbid L^2E^c, $LQ\tilde{D}^c$, and LH_u, we require $\eta_L \neq \eta_{H_d}$, and to forbid $\tilde{D}^2\tilde{U}^c$, we require $\eta_{U^c} \neq (\eta_{D^c})^{-2}$ or $\eta_{\tilde{D}^c}^2\eta_{\tilde{U}^c} = \eta_Q^{-3}\eta_{H_u} \neq 1$. If, as in the FIQS model, the Q_I all have the same $U(1)$ charges, the constraint that they have the same R charge implies that $F^I - F^J = 2ni\pi$. Then, since we also require $\sum_I F^I = 2mi\pi$, it is easy to check that $F^I = 2n^Ii\pi$, giving $\eta_{MJ} = e^{2i\pi\sum_I n^I[(1/3)\zeta^M-\zeta_I^M]}$, $\eta_M = e^{2i\pi\sum_I n^I[(1/3)(\zeta^M-1)-\zeta_I^M]}$. We can find an R parity provided there is some compactification for which we can identify the particles of the MSSM in such a way that the above constraints are satisfied. With the above choices [$Q_I, T^c, H_u \in U; \tilde{U}^c, \tilde{D}^c, L, E^c, H_d \in T$ and degenerate $U(1)$ charges for fixed flavor in each sector], they take the form

$$\eta_Q = e^{2i\pi\beta}\eta_{H_d}^{-2}, \qquad \eta_{\tilde{U}^c} = \eta_{T^c} = e^{-2i\pi\beta}\eta_{H_d}^3, \qquad \eta_{\tilde{D}^c} = e^{-2i\pi\beta}\eta_{H_d}, \qquad \beta \neq \frac{n}{3}, \qquad \eta_{H_u} = \eta_{H_d}^{-1},$$
$$\eta_L = e^{2i\pi\alpha}\eta_{H_d}, \qquad \eta_{E^c} = e^{-2i\pi\alpha}\eta_{H_d}^{-2}, \qquad 0 < \alpha, \beta < 1. \tag{14}$$

When the electroweak symmetry is broken, we redefine R parity as before, so that $\eta_{H_d} = 1$.

Other scenarios can be considered. For example, a $Q_3 T^c H_u$ coupling is allowed if these are all in the twisted sector T, their $U(1)$ charges sum to zero, and there is no Π factor in (7) required by a further string symmetry. In this case, it would be possible to have all quarks of the same flavor having the same $U(1)$ charge. It is in fact not necessary to have identical $U(1)$ charges to assure equal masses for squarks and sleptons of the same flavor, which is what is actually needed to avoid unwanted FCNC; scalars ϕ^M with the same value of ζ^M have the same masses, but could have different values of ζ_I^M. (In the FIQS model considered above, the MSSM candidates with the same flavor that are degenerate in ζ^M are also degenerate in ζ_I^M.) Models with vanishing or very small values of ζ^M for MSSM particles are favored by experimental and theoretical constraints if symmetry is broken by gaugino condensation, so the values of the ζ_I^M could be the governing factors in determining R parity.

To what extent have we achieved the conventional definition of R parity? With appropriate choices of phases in (14) we achieve the elimination of baryon (B) and lepton (L) number violating couplings of dimension two or three in the superpotential. [Fast proton decay is avoided by eliminating either one of these, but the $\tilde{U}^c(\tilde{D}^c)^2$ coupling by itself would induce neutron-antineutron oscillations.] Higher dimension operators can generate B and L violation, as is the case with the conventional definition of R parity. In the latter case, the R-allowed dimension-four operator $\tilde{U}^c\tilde{U}^c\tilde{D}^c E^c$ in the superpotential leads to dimension-five operators in the effective Lagrangian that may be problematic [15] even if these couplings are Planck- or string-scale suppressed, given the current bounds on the proton decay lifetime. This problem is easily evaded in the current context provided $3\beta + \alpha \neq n$. The stability of the lightest neutralino is assured at the same level as proton stability since its decay products would have to include an odd number of standard model fermions and hence violate B and/or L.

In conclusion, effective supergravity from the weakly coupled heterotic string may possess a residual discrete symmetry that plays the role of R parity if the T moduli are stabilized at self-dual points when supersymmetry is broken by condensation in a hidden sector. For example, the requirement that a Z_3 orbifold compactification accommodate R parity imposes constraints on the quantum numbers of chiral fields. By way of illustration, a specific model that fails this test was analyzed.

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