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A NEW APPROACH TO SUPERSYMMETRY BREAKING IN SUPERSTRING MODELS

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ABSTRACT

We re-examine supersymmetry breaking in effective four-dimensional field theories derived from the superstring, noting three features of modern four-dimensional string theories that were absent in the earlier compactifications of ten-dimensional string theories. These features are: the replacement of gauge-singlet moduli fields by gauge non-singlet chiral supermultiplets, vector-like matter fields in the hidden sector, and extra $U(1)$ gauge symmetries under which both hidden and observable matter fields transform. We point out that an anomalous $U(1)$ gauge symmetry can fix the magnitudes of the *ersatz* moduli fields. We propose a new scenario in which composite chiral supermultiplets in the hidden sector break global supersymmetry, and the extra massive $U(1)$ gauge supermultiplets communicate this breaking to the observable sector via radiative corrections in the renormalizable effective field theory.

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One of the greatest obstacles to the derivation of realistic four-dimensional field theories from the superstring has been our lack of understanding of the dynamical determination of scales, in particular that of supersymmetry breaking. The practical wish to use supersymmetry to protect the gauge hierarchy argues for a small scale of global supersymmetry breaking $\sim 1\ TeV$ in the observable sector [1]. However, cosmological arguments suggest that the scale of local superymmetry breaking should be much larger [2]. The only known mechanisms for supersymmetry breaking in string perturbation theory relate its magnitude to internal radii, and therefore give very large sparticle masses $\sim m_P$ [3]. Moreover, they break supersymmetry locally and globally at the same time and by the same amount, so that gaugino and gravitino masses are identical and large. Therefore, *faute de mieux*, phenomenological attention has focussed on non-perturbative sources of supersymmetry breaking in effective four-dimensional field models obtainable from string theories. Historically the first mechanism of this type to be proposed was condensation of gauginos in a hidden gauge group sector [4,5], possibly accompanied by condensation of a higher-dimensional antisymmetric tensor field [5]. Subsequently, the possibility was raised [6] that an anomalous four-dimensional $U(1)$ gauge symmetry might break supersymmetry *à la* Fayet-Iliopoulos [7], and gravitino condensation was proposed [8] as an alternative to antisymmetric tensor condensation.

None of these mechanisms is completely satisfactory. It is not even clear whether gaugino condensation occurs in all hidden sector gauge groups [9]. If it does, it is not clear how the scale of supersymmetry breaking could be much smaller than m_P, since antisymmetric tensor [10] or gravitino condensates [8] would occur at the Planck scale. It has also never become clear how the supersymmetry breaking in the hidden sector is fed through to the observable sector, though there exist various suggestions why the observable sector supersymmetry breaking scale may be suppressed by radiative corrections relative to that in the hidden sector [11].

Most analyses of the gaugino condensation mechanisms [4, 5] have used the specific form of low-energy effective four-dimensional supergravity theory derived from a toy model for Calabi-Yau compactification of the ten-dimensional heterotic string with one observable matter generation [12]. Since then, more realistic four-dimensional string models have been constructed [13], whose effective field theories [14-16] have certain similarities to the original toy model, but also many differences that could play key rôles. Among these should be mentioned the replacement

of a single gauge singlet modulus field by a set of gauge non-singlet matter fields [14,15], the presence of vector-like matter fields in the hidden sector [16], and the existence of extra $U(1)$ gauge symmetries under which both hidden and observable sector matter fields transform [16].

In this paper we formulate a new scenario for the dynamical determination of scales, including that of supersymmetry breaking, in string theories, by exploiting these new features. We first point out that an anomalous $U(1)$ gauge symmetry can be used to fix the analogues of the moduli fields related to "compactification radii". We then point out that when coupled to an appropriate no-scale supergravity theory [1] , the non-perturbative superpotential for composite vector and matter supermultiplets [17] can induce a breakdown of global supersymmetry for both vector and matter particles in the hidden sector. The matter supersymmetry breaking can then be communicated to the observable sector by extra massive $U(1)$ gauge supermultiplets via radiative corrections calculated in the renormalizable effective low-energy four-dimensional field theory. The magnitude of supersymmetry breaking in the observable sector can easily be much lower than the scale at which the hidden sector gauge interactions become strong and the masses of the extra $U(1)$ gauge bosons, and may easily be $0(1)$ TeV as desired to make the gauge hierarchy natural.

We start by recapitulating the form [12] of effective supergravity theory that has previously been assumed in most analyses of supersymmetry breaking in superstring-inspired field theories [11], and point out the different features that arise in realistic four-dimensional string models [14,15], illustrating them with a specific flipped $SU(5) \times U(1)$ string theory [16]. The Kähler potential has often been assumed to take the form [11]

$$G = -\ell n(S + S^+) - 3\ell n(T + T^+ - \sum_i |\phi_i|^2) \tag{1}$$

where S and T are gauge-singlet chiral superfields whose spin-zero components are combinations of the ten-dimensional supergravity dilaton, a "compacton" scaling six compactified dimensions of an assumed underlying ten-dimensional heterotic string, and two axions associated with components of the ten-dimensional antisymmetric tensor field, and the ϕ_i are gauge non-singlet chiral superfields. The gauge kinetic function is moreover taken to be

$$f_{ab} = \delta_{ab}\ .\ S \tag{2}$$

Within this class of models, although a non-perturbative mechanism to fix the S field has been proposed [5], no one has ever figured out how to determine the T field which corresponds to the moduli of a "realistic" compactification.

However, in general there is no need to invoke an ancestral compactified ten-dimensional heterotic string, and larger gauge symmetries can be realized by formulating string theories directly in four dimensions [13]. These can be visualized partially as corresponding to special choices of the magnitudes of compactification moduli at which the effective four-dimensional gauge symmetry is enhanced. In all four-dimensional models, the S field appears as a $d = 4$ dilaton, and the $-\ell n(S+S^+)$ term in the Kähler potential (1) and the gauge kinetic function (2) are unchanged [14,15]. However, the T field is in general replaced by a set of gauge non-singlet fields, and the form of the corresponding part of the Kähler potential is modified, though it always takes a generalized no-scale form. For example, in four-dimensional string models formulated using world-sheet fermions with periodic and antiperiodic boundary conditions, which are related to Z_2 orbifolds, the corresponding piece of the Kähler potential is [14,15]

$$G = -\ell n(S + S^+) - \sum_{a=1}^{3} \ell n \; \eta_a^0 + \dots :$$
$$\eta_a^0 \equiv \frac{1}{\sqrt{2}} \left| 1 + |\hat{\eta}_a^1|^2 - \sum_{i_a=2}^{n_a} |\hat{\eta}_a^{i_a}|^2 \right|, \qquad (3)$$
$$\hat{\eta}_a^{1^2} \equiv -1 + \sum_{i_a=2}^{n_a} \hat{\eta}_a^{i_a 2}$$

In such a model, the value of e^G is determined by the expectation values of S and of the gauge non-singlet fields $\hat{\eta}_a^{i_a}$. Unlike the previous T field, the latter appear in D-terms and possibly a non-trivial superpotential coupled to matter fields, and so are subject to the restrictions of D- and F-flatness, which constrain their vacuum expectation values.

It has normally been assumed [11] that the four-dimensional gauge group is the product of an observable sector group and a non-Abelian hidden sector group. It has further been assumed [11] that the chiral superfields transform non-trivially only under the observable sector gauge group. Four-dimensional string theories usually also contain a hidden sector gauge group, but

may also contain matter superfields transforming non-trivially under the non-Abelian hidden gauge group. An example is the model of Ref. [16] whose observable sector gauge group is $SU(5) \times U(1)$ and whose hidden sector gauge group is $SO(10) \times SO(6)$. The spectrum of light matter fields include five $\underline{10}$ representations of $SO(10)$, and five $\underline{6}$ and six $\underline{4}+\bar{\underline{4}}$ representations of $SO(6)$. Note that the hidden matter fields belong in this case to a self-conjugate representation of the hidden sector gauge group, which is a necessary feature of realistic models.

Finally, we note that both hidden sector matter and the observable matter fields may transform non-trivially under extra $U(1)$ gauge group factors. Specifically, in the $SU(5) \times U(1)$ model [16] mentioned above, there are four extra $U(1)$ gauge symmetries, all of which are spontaneously broken at some large scale. Thus there are gauge couplings between hidden and observable sector particles, although these would be of very short range and hence very weak if the corresponding $U(1)$ vector bosons are very massive as we expect.

Having introduced these novel possible features of low-energy four-dimensional field theories derived from string, we will now indicate how they permit promising new mechanisms for fixing scales, in particular that of supersymmetry breaking.

In general, one of the $U(1)$ gauge symmetries is anomalous [6], and this is certainly the case in the $SU(5) \times U(1)$ model [16] introduced above. The D-term of an anomalous $U(1)_A$ is:

$$D_A = \frac{g}{192\pi^2} Tr \; Y_A + \sum_i Y_A^i |\phi^i|^2 \qquad (4)$$

where $g^2 = \frac{1}{2}/(S+S^+)$ is the gauge coupling, and the Y_A^i are the anomalous $U(1)_A$ charges of the matter fields ϕ^i. It is in general possible [18] to avoid catastrophic large supersymmetry breaking accompanied by a large cosmological constant if there are large v.e.v.'s $< 0|\phi^i|0 >= 0(\sqrt{g})$. For example, in the $SU(5) \times U(1)$ model one such desirable pattern of v.e.v.'s giving vanishing D-terms for all the gauge symmetries has [16]

$$< 0||\hat{\eta}_1|^2|0 >= \frac{1}{2} < 0||\hat{\eta}_2|^2|0 >= \frac{180}{15.192\pi^2} g, \quad < 0||\hat{\eta}_3|^2|0 >= 0 \qquad (5)$$

where $|\hat{\eta}_1|^2 = |\bar{\Phi}_{23}|^2$, $|\hat{\eta}_2|^2 = |\Phi_{31}|^2$ and $|\hat{\eta}_3|^2 = |\Phi_{12}|^2$ in the notation of Ref. [16]. Notice that with $g = 0(1)$ as required by experiment*, these v.e.v.'s are much smaller than the Planck scale, thereby justifying a perturbative treatment. The v.e.v.'s (5) are also consistent with the F-flatness requirement imposed by $\bar{\Phi}_{23}\bar{\Phi}_{31}\bar{\Phi}_{12}+$ herm. conj. terms in the superpotential [16]. Inserting the v.e.v.'s (5) into the expression (3) for the Kähler potential, we see that $\eta_a^0 \simeq \sqrt{2}$ for $a = 1,2,3$, and hence the three analogues of the conventional moduli fields have been determined to be of order unity. We consider that this mechanism is likely to be of general applicability. As was remarked earlier, in generic four-dimensional string models the gauge symmetry is enhanced for some specific values of the moduli. When the anomaly breaks this gauge symmetry perturbatively, the moduli are fixed at values close to these specific ones.

Next, we analyze the strongly-coupled non-Abelian hidden sector gauge dynamics, following the analysis in Ref. [17] for globally supersymmetric models. We introduce a composite supermultiplet $H \equiv (W_a W^a)^{\frac{1}{3}}$ (where W_a is the conventional gauge supermultiplet, with a an adjoint index), and composite matter supermultiplets $U^{ij} \equiv Tr(\phi^i\phi^j)$ (where the ϕ^i are different hidden sector matter supermultiplets, and the trace runs over the matter representation index), but for simplicity we will drop the index i†. As was shown in Ref. [17], the effective superpotential for the "pion-like" composite supermultiplets takes the general form

$$W(H,U) = de^{-\lambda S}H^3\left[\ell n\left(\frac{H^pU^q}{\Lambda_H^{p+2q}}\right) - \frac{p}{3}\right] + c$$
$$\equiv de^{-\lambda S}H^3\left[L - \frac{p}{3}\right] + c \qquad (6)$$

where the value of λ depends on the hidden sector-gauge group, d is a non-perturbatively determined number, the powers p and q are determined by global $U(1)$ anomalies‡, and we have added a constant term that could be generated by antisymmetric tensor [5] or gravitino condensation [8]. In the case of a globally-supersymmetric theory [17], the factor $e^{-\lambda S}$ and the constant c would be absent. We ignore the possible existence of a "current" mass term for the

* The value of g in this scenario is determined by that of $< 0|S + S^+|0 >$, which is fixed by the non-perturbative effects discussed below.

† In the $SU(5) \times U(1) \times SO(10) \times SO(6)$ model [16], we expect the hidden $SO(10)$ group to be the most important for supersymmetry breaking, and we have $i = 1,...,5$ for the $\underline{10}$ matter representations.

‡ For an $SO(10)$ hidden sector with $N\underline{10}$ matter representations, we have $p = 3(8 - N)$ and $q = 1$.

matter superfield, which could appear due to non-renormalizable interactions in some of the class of superstring-derived models that we consider, but would introduce no essential change in the scenario we now outline. We do not determine the form of the effective Kähler metric for the composite superfields H and U, but assume for simplicity that it takes the same form as that found earlier [14,15] for elementary matter superfields, which takes a simple form [14,15] when the moduli η_a^0 are close to their symmetry values

$$G_{H,U} \simeq -\ell n(1 - |H|^2) - \ell n(1 - |U|^2) \qquad (7)$$

The working of our scenario for supersymmetry breaking is not sensitive to the details of $G_{H,U}$ as we will see when we analyze the effective potential given by the Kähler potential $G = -\ell n(S + S^+) + G_{H,U} + \ell n|W(H,U)|^2$.

The effective scalar potential we study is

$$V = e^G \sum_{i=S,H,U} G_i(G_{ii*})^{-1}G_i^+ \qquad (8)$$

where we have already taken into account the cancellation of the $(-3e^G)$ term by the η_a^0-dependent pieces from (3), and we need no longer consider the variation of the η_a^0 since they are fixed by the $U(1)_A$ D-term (4). Using Eqs. (3), (6) and (7) to evaluate (8), we find that

$$V \simeq \frac{1}{2^3(S+S^+)}\left(\frac{1}{1-|H|^2}\right)\left(\frac{1}{1-|U|^2}\right)\times$$
$$\times\left\{|W + \lambda de^{-\lambda S}(S+S^+)H^3(L - \frac{p}{3})|^2 + |H^*W + 3(1-|H|^2)de^{-\lambda S}H^2L|^2\right.$$
$$\left. + |U^*W + (1-|U|^2)dqe^{-\lambda S}H^3/U|^2\right. \qquad (9)$$

In the past, most authors [11] have ignored the possible presence of the U field, have assumed that $< 0|H|0 >$ is dynamically determined, and then used the analogue of the first, $|G_S|^2$, term in (8) to fix $< 0|S + S^+|0 >$ in such a way that $V = 0$ whilst local and global supersymmetries are broken: $W \neq 0$ and $W_S \neq 0$. This conventional approach was recently extended in the last paper of Ref. [11] to include the H field explicitly, justifying the earlier approach. We now go further, showing that there exists a global and local supersymmetry-breaking vacuum with $W, W_S, W_H, W_U \neq 0$, which will then have novel consequences.

We assume that $\Lambda_H << 0(1)$ and look for solutions in which $h \equiv |H|$ is small: $h \equiv \Lambda_H^\alpha$: $\alpha > 0$, in which case the precise form of H-dependence in (7) is not important. The $|G_H|^2$ and $|G_U|^2$ terms in (9) vanish simultaneously if

$$\frac{h^2}{1-h^2} = \frac{3}{q} \quad \ell n\left(\frac{h^p u^q}{\Lambda_H^{p+2q}}\right) \quad \frac{u^2}{1-u^2} \tag{10}$$

where $u = |U|$. When both h and u are small, equation (10) may be rewritten as

$$\frac{h^p u^q}{\Lambda_H^{p+2q}} = \exp\left(\frac{q}{3} h^2/u^2\right) \tag{11}$$

Clearly, when $1 + 2q/p > \alpha > \frac{1+2q/p}{1+q/p}$ we have

$$h = \Lambda_H^\alpha, \quad u \simeq \Lambda_H^{2+(1-\alpha)p/q} << 1 \qquad \left(1 + 2q/p > \alpha > \frac{1+2q/p}{1+q/p}\right) \tag{12}$$

More analysis is required when $\alpha \geq 1 + 2q/p$, in which case $u = 0(1)$ and one becomes sensitive to the detailed form of $G_{H,U}$ (7), and there may even no longer exist a non-trivial $V = 0$ vacuum. However, we will be interested in small α : $\frac{1+2q/p}{1+q/p} > \alpha > 0$, in which case the solution to (11) is

$$h = \Lambda_H^\alpha, \quad u \simeq \frac{\Lambda_H^\alpha}{\sqrt{\frac{3}{q}[p(\alpha-1) + q(\alpha-2)]\ell n \Lambda_H}} \tag{13}$$

The $|G_S|^2$ and $|G_H|^2$ terms in (9) vanish simultaneously if

$$(S + S^+) = \frac{3}{\lambda h^2} \cdot \frac{L}{L - p/3} \tag{14}$$

In the case where $\alpha > \frac{1+2q/p}{1+q/p}$ (12), we would have

$$L \simeq \frac{q}{3}\frac{h^2}{u^2} \simeq \frac{q}{3}\Lambda_H^{2(1+p/q)\alpha - 2(2+p/q)} << 1 \tag{15}$$

which gives no solutions to (14) with $(S + S^+) = 2/g^2 > 0$ as required. However, when

$\frac{1+2q/p}{1+q/p} > \alpha > 0$ (13), Eq. (14) is solved by

$$S + S^+ \simeq \frac{3}{\lambda \Lambda^{2\alpha}} >> 0 \qquad \left(\frac{1+2q/p}{1+q/p} > \alpha > 0\right) \tag{16}$$

as required. To complete our analysis, we must now see how the *a priori* unknown parameter α is determined by requiring $G_S = 0$:

$$c + (1 + \lambda(S + S^+))de^{-\lambda S} H^3 (L - p/3) = 0 \tag{17}$$

Substituting the previous expressions (13), (16) for H, U and S into Eq. (17), we find

$$\ell n\left(\frac{-c}{d}\right) \simeq -\frac{3}{2\Lambda_H^{2\alpha}} + \alpha \, \ell n \, \Lambda_H + ... \tag{18}$$

In other words, it is the ratio of the as-yet-unknown non-perturbatively-determined numbers c, d that fixes α and hence all the v.e.v.'s.

The $V = 0$ vacuum that we have found is generic, and our procedure for finding it can be rephrased as an expansion in the small parameter $\frac{g^2}{4\pi}$, which is justified if the ratio of non-perturbative numbers $|\frac{c}{d}| << 1$. As was mentioned above, in this vacuum both local and global supersymmetries are broken, since $W, W_S, W_H, W_U \neq 0$. In the past, the fact that W and $W_S \neq 0$ has been used [11] to seed global supersymmetry breaking via mass splittings in the graviton/gravitino and S supermultiplets. This was then thought to be fed through to the observable-sector matter fields by loop corrections. However, there are no renormalizable field theory interactions coupling the graviton/gravitino, S or H supermultiplets to the observable sector fields, and hence these putative loop corrections could not be calculated reliably. However, now the situation is different, because there is also global supersymmetry breaking for the U supermultiplet, which has in general extra $U(1)$ gauge interactions coupling it to the observable sector.

We do not have space in this paper to present detailed calculations of the renormalizable loop corrections, which would in any case be somewhat model-dependent, but we point out a few key features. Figure 1 is a superfield diagram coupling the U supermultiplet with broken global supersymmetry indicated by a cross, to a communicating extra $U(1)$ gauge supermultiplet $V^{(1)}$. Since the common extra $U(1)$ gauge supermultiplet mass $m_{V^{(1)}}$ is expected to be proportional to the v.e.v.'s of the moduli fields (5), and hence large, the generic magnitude of supersymmetry breaking in the $V^{(1)}$ supermultiplet will be [19]

$$\tilde{m}_{V^{(1)}} \sim \Lambda_H (\frac{\Lambda_H}{m_P})^m \ : \ m > 0 \tag{19}$$

Figure 2 is a superfield diagram coupling the $V^{(1)}$ supermultiplet to a generic observable-sector matter supermultiplet O^i, where again we expect that since $\tilde{m}_{V^{(1)}} << m_{V^{(1)}}$:

$$\tilde{m}_{O^i} \sim \Lambda_H (\frac{\Lambda_H}{m_P})^{m'} \ : \ m' > m > 0 \tag{20}$$

Finally, Fig. 3 is a renormalizable superfield diagram coupling the O^i to a generic observable sector gauge supermultiplet $V^{(0)}$, for which we expect supersymmetry breaking

$$\tilde{m}_{V^{(0)}} \sim (\frac{\alpha}{\pi}) \tilde{m}_{O^i} f(\frac{\tilde{m}_{O^i}}{m_{V^{(0)}}}) \tag{21}$$

Because of the large communicating gauge boson mass, this scenario offers a strong suppression of the renormalizable loop contributions* to supersymmetry breaking in the observable sector.

We have seen in this paper how three novel features of realistic four-dimensional string models offer a new scenario for fixing scales, in particular that of supersymmetry breaking. The axial $U(1)$ gauge anomaly determines the analogues of the compactification moduli. Gauge non-singlet hidden-sector matter fields break both local and global supersymmetries. Extra massive $U(1)$ gauge supermultiplets then communicate this symmetry breaking to the observable sector with a greatly reduced magnitude. We have illustrated this scenario with an $SU(5) \times U(1) \times SO(10) \times SO(6)$ superstring model proposed recently [16]. We plan to present soon a more detailed exploration of this scenario in the context of that specific model.

* There may still be hard-to-evaluate non- renormalizable loop contributions as well.

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FIGURE CAPTIONS

1) Superfield diagram communicating supersymmetry breaking (marked by ⊗) from the composite hidden sector chiral superfield U to the massive vector supermultiplet $V^{(1)}$.

2) Superfield diagram communicating supersymmetry breaking (marked by ⊗) from the massive vector supermultiplet $V^{(1)}$ to the light observable-sector matter fields $O^{(0)}$.

3) Superfield diagram communicating supersymmetry breaking (marked by ⊗) from the observable-sector matter fields $O^{(0)}$ to the light observable-sector gauge fields $V^{(0)}$.

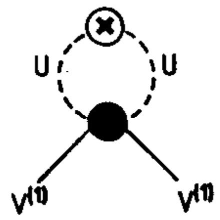

Fig. 1

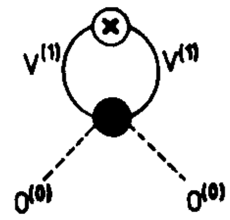

Fig. 2

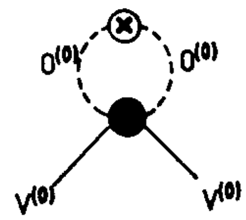

Fig. 3