Minimal $SO(10)$ Unification*

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Abstract

It is shown that the doublet-triplet splitting problem can be solved in $SO(10)$ using the Dimopoulos-Wilczek mechanism with a very economical Higgs content and simple structure. Only one adjoint Higgs field is required, together with spinor and vector fields. The successful SUSY GUT prediction of gauge coupling unification is preserved. Higgsino-mediated proton decay can be suppressed below (but not far below) present limits without fine-tuning.

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1 Introduction

The striking unification of gauge couplings\textsuperscript{1} at about $10^{16}$ GeV in the minimal supersymmetric standard model (MSSM) points toward the possibility of a supersymmetric grand unified theory (SUSY GUT). The minimal SUSY $SU(5)$ prediction of $\sin^2 \theta_W$ is $0.2334 \pm 0.0036$, while the experimental value is $0.2324 \pm 0.0003$.\textsuperscript{1}

$SO(10)$ is generally thought to be the most attractive grand unified group for a number of reasons. It achieves complete quark-lepton unification for each family, explains the existence of right-handed neutrinos and of “see-saw” neutrino masses, has certain advantages for baryogenesis, in particular since $B - L$ is broken\textsuperscript{2}, and has the greatest promise for explaining the pattern of quark and lepton masses.\textsuperscript{3}

The greatest theoretical problem that any grand unified theory must face is the gauge hierarchy problem,\textsuperscript{4} and in particular that aspect of it called the “doublet-triplet splitting problem”\textsuperscript{5}, or 2/3 splitting problem for short. The only way to achieve natural 2/3 splitting in $SO(10)$ is the Dimopoulos-Wilczek (DW) mechanism.\textsuperscript{6} In Ref. 7 it was shown that realistic $SO(10)$ models can be constructed using the DW mechanism.

One criticism that is sometimes made about $SO(10)$ is that solving the 2/3 splitting problem requires a somewhat complicated Higgs structure. The models constructed in Ref. 7 and 8 contained at least the following Higgs multiplets: three adjoints ($45$), two rank-two symmetric tensors ($54$), a pair of spinors ($16 + 16$), two vectors ($10$), and several singlets. Aside from the issue of simplicity, there are some indications that it is difficult to construct grand unified models with a multiplicity of adjoint fields from superstring theory.\textsuperscript{9}

The necessity of a complicated Higgs structure is largely traceable to one technical problem, namely the breaking of the rank of $SO(10)$ from five to four without destabilizing the DW solution of the gauge hierarchy problem. The complete breaking of $SO(10)$ to the Standard Model requires at least two sectors of Higgs: an adjoint sector and a spinor sector. The adjoint sector plays the central role in the DW mechanism for 2/3 splitting, while the spinor sector both breaks the rank of $SO(10)$ and gives right-handed neutrinos mass.

The dilemma is that if the spinor sector is coupled to the adjoint sector it tends to destabilize the DW form of the adjoint vacuum expectation value.
required for the 2/3 splitting, while if the two sectors are not coupled (or coupled very weakly\textsuperscript{10,11}) to each other in the superpotential there arise colored and charged (pseudo)goldstone fields due to the fact that that certain generators of $SO(10)$ are broken by both the adjoint and the spinors, and in particular the generators that transform as $(3,2,\frac{1}{6}), (\overline{3},2,-\frac{1}{6}), (\overline{3},1,-\frac{2}{3})$, and $(3,1,\frac{2}{3})$. These destroy the unification of couplings\textsuperscript{10}.

In Refs. 7 and 8, an indirect way to couple the two sectors together without destabilizing the hierarchy was found. However, this solution to the problem involved a somewhat complicated Higgs structure including at least three adjoint fields.

In this letter we show that there is a very simple way to couple the spinor and adjoint sectors together, with a stable hierarchy and no pseudo-goldstones. Before describing it, it will be helpful to review in more detail the problems that have been discussed above.

2 The problem of a stable hierarchy in $SO(10)$

The DW mechanism is based on the existence of an adjoint Higgs field, which we shall call $A$, getting a vacuum expectation value (VEV) in the $B - L$ direction:

$$\langle A \rangle = \begin{pmatrix}
0 \\
0 \\
a \\
\bar{a}
\end{pmatrix} \otimes i\tau_2, \quad (1)$$

where $a \sim M_G$, the unification scale. The lower-right $3 \times 3$ block corresponds to $SU(3)$ of color, and the upper-left $2 \times 2$ block to Weak $SU(2)$. When this adjoint is coupled to vector representations, which we will call $T_1$ and $T_2$, by terms such as $T_1 A T_2$ the color-triplets in the vectors are given superlarge masses, while the Weak doublets remain massless. (By having also a term of the form $M_T T_2^T$ it is ensured that only one pair of Weak doublets remains light. The proton-decay amplitude from the exchange of colored Higgsinos is proportional to $M_T/a^2$, so that if $M_T/a \lesssim 10^{-1}$ it is sufficiently suppressed\textsuperscript{7,8}) Such a “DW form” for the adjoint VEV is not possible in $SU(5)$, where $\text{tr}(A) = 0$. 
An adjoint alone is not sufficient to break $SO(10)$ to the Standard Model group, $G_{SM}$, and in particular cannot reduce the rank of the group. This requires either spinorial Higgs $(16 + \overline{16})$ or rank-five antisymmetric tensor Higgs $(126 + \overline{126})$. As the latter tend to destroy the perturbativity of the unified interactions below the Planck scale, we will assume that the rank-breaking sector has spinors, which we shall call $C$ and $\overline{C}$. These spinors (also necessary to give mass to the right-handed neutrinos) have VEVs in the $SU(5)$-singlet direction.

The spinor VEVs break $SO(10)$ down to $SU(5)$, and thus the sector of the superpotential which depends on $C$ and $\overline{C}$ but not on $A$, which we shall call $W_C$, leaves massless at least those components of the spinors in the coset $SO(10)/SU(5)$. That is just a $10 + \overline{10} + 1$ of $SU(5)$, or a $[(3, 2, \frac{1}{6}) + (\overline{3}, 1, -\frac{2}{3}) + (1, 1, +1) + H.c.] + (1, 1, 0)$ of $G_{SM}$.

The adjoint $A$, with the VEV shown in Eq.(1), breaks $SO(10)$ down to $SU(3)_c \times U(1)_{B-L} \times SO(4)$. ($SO(4) = SU(2)_L \times SU(2)_R$.) The part of the superpotential that depends on $A$ but not on the spinors, which we shall call $W_A$, thus leaves massless at least those components of the adjoint in the cosets $SO(10)/(SO(6) \times SO(4))$ and $SO(6)/(SU(3)_c \times U(1)_{B-L})$. The first of these cosets consists of a $(6, 4)$ of $SO(6) \times SO(4)$, which contains $[(3, 2, \frac{1}{6}) + (3, 2 - \frac{5}{6}) + H.c.]$ of $G_{SM}$. The second coset consists of $[(\overline{3}, 1, -\frac{2}{3}) + H.c.]$ of $G_{SM}$.

Thus, with no coupling between the two sectors, there are extra, uneaten goldstone fields in $(3, 2, \frac{1}{6}) + (\overline{3}, 1, -\frac{2}{3}) + H.c.$ To avoid these, the adjoint must couple to the spinor. The obvious way to couple them together, by the term $g\overline{C} A C$, directly destabilizes the DW form assumed for $\langle A \rangle$. Let us assume that $A$ has the form diag$(b, b, a, a, a) \otimes i\tau_2$. $b = 0$ is the desired DW form. The VEVs of the spinors point in the $SU(5)$-singlet direction and have magnitude $c_0 \sim M_{GUT}$. Then $g\overline{C} A C = -\frac{g}{2}(2b + 3a)c_0^2$. The terms $W_A$ must have a form that gives $\partial W_A/\partial b = O(M_{GUT})b$, so that by themselves they would give $b = 0$. But taking into account also the coupling term $g\overline{C} A C$, one has $0 = -F_b^* = \partial W_{tot}/\partial b = O(M_{GUT})b - gc_0^2$, or $b \sim gM_{GUT}$. In the DW mechanism the Higgs doublets get a “see-saw” mass of order $b^2/M_{GUT}$, so that $g$ must be less than about $10^{-7}$. This is the assumption made in Ref. 11. This leads to the pseudo-goldstone fields getting masses only of order $gM_{GUT} < 10^9$ GeV, and thus to $\sin^2 \theta_W = 0.2415$.

The spinor sector and adjoint sector must be coupled together in some more subtle way. In Ref. 7 such a way was proposed. There the spinor sector
was assumed to contain a different adjoint Higgs, called $A'$, whose VEV points in the SU(5) singlet direction. The two sectors (namely the $A$ sector and the $(C, \overline{C}, A')$ sector) were then coupled together by a term $\text{tr}AA'A''$, where $A''$ was a third adjoint. Because the $AA'A''$ is totally antisymmetric under the interchange of any two adjoints (due to the fact that the adjoint is an antisymmetric tensor), there have to be three distinct adjoints in this term. This antisymmetry ensures, as it is easy to see, that this term does not contribute to any of the $F$ terms as long as the VEVs of the three adjoints commute with each other. Therefore it does not destabilize the DW form of the VEV of $A$. And yet, it can also be shown that this trilinear term is sufficient to prevent the existence of any pseudo-goldstone fields.

This has been the choice until now: to assume a complicated Higgs sector with at least three adjoint Higgs fields or to assume that light pseudo-goldstones exist which disturb the beautiful SUSY GUT prediction of the Weak angle.

3 Solving the problem

The solution to the above difficulty turns out to be remarkably simple. Let there be a single adjoint field, $A$, and two pairs of spinors, $C + \overline{C}$ and $C' + \overline{C}'$. The complete Higgs superpotential is assumed to have the form

$$W = W_A + W_C + W_{ACC'} + (T_1AT_2 + ST_2^2).$$

(2)

The precise forms of $W_A$ and $W_C$ do not matter, as long as $W_A$ gives $\langle A \rangle$ the DW form, and $W_C$ makes the VEVs of $C$ and $\overline{C}$ point in the SU(5)-singlet direction. For specificity we will take $W_A = \frac{1}{4M} \text{tr}A^4 + \frac{1}{2}P_A(\text{tr}A^2 + M_A^2) + f(P_A)$, where $P_A$ is a singlet, $f$ is an arbitrary polynomial, and $M \sim M_G$. (It would be possible, also, to have simply $m\text{tr}A^2$, instead of the two terms containing $P_A$. However, explicit mass terms for adjoint fields may be difficult to obtain in string theory.) We take $W_C = X(\overline{C}C - P_C^2)$, where $X$ and $P_C$ are singlets, and $\langle P_C \rangle \sim M_G$.

The crucial term that couples the spinor and adjoint sectors together has the form

$$W_{ACC'} = C' \left( \left( \frac{P}{M_P} \right) A + Z \right) C + \overline{C}' \left( \left( \frac{P}{M_P} \right) A + \overline{Z} \right) C',$$

(3)
where $Z$, $\overline{Z}$, $P$, and $\overline{P}$ are singlets. $\langle P \rangle$ and $\langle \overline{P} \rangle$ are assumed to be of order $M_G$. The critical point is that the VEVs of the primed spinor fields will vanish, and therefore the terms in Eq. (3) will not make a destabilizing contribution to $-F_A^* = \partial W/\partial A$. This is the essence of the mechanism.

$W$ contains several singlets ($P_C$, $P$, $\overline{P}$, and $S$) that are supposed to acquire VEVs of order $M_G$, but which are left undetermined at tree-level by the terms so far written down. These VEVs may arise radiatively when SUSY breaks, or may be fixed at tree level by additional terms in $W$, possible forms for which will be discussed below.

The VEVs of $A$ and $P_A$ are determined by the equations $0 = -F_A^* = \frac{1}{M} A^3 + P_A A$ and $0 = -F_{P_A}^* = \frac{1}{2} (\text{tr} A^2 + M_A^2) + f'(P_A)$. If $\langle A \rangle = \text{diag}(a_1, a_2, a_3, a_4, a_5) \otimes i\tau_2$, then the first equation implies that $a_i^2 = 0$ or $M \langle P_A \rangle \equiv a^2$, for each $i$. There is, therefore, a discrete vacuum degeneracy. The DW vacuum is obtained if two of the $a_i$'s vanish and the other three have the same sign and magnitude $a$. In that case, $\text{tr} A^2 = -6M \langle P_A \rangle$ and $\langle P_A \rangle$ is determined by $0 = f'(\langle P_A \rangle) - 3M \langle P_A \rangle + M_A^2/2$.

$F_X = 0$ implies that $\langle CC \rangle = \langle P_C \rangle^2 \sim M_G^2$. The $D$ terms and soft, SUSY-breaking terms will ensure that $\langle C \rangle \equiv \langle C \rangle \sim M_G$.

The most interesting equations are

$$0 = -F_{C'}^* = \left( \frac{P}{M_P} \right) A + Z \right) \overline{C}, \quad (4)$$

and

$$0 = -F_{\overline{C'}}^* = \overline{C} \left( \frac{P}{M_P} \right) \left( A + Z \right). \quad (5)$$

It is only necessary to consider the first of these two equations, as they have the same structure. Let the VEV of $C$ be decomposed as follows: $\langle C \rangle = \sum_K f_K C_K$, where the $C_K$ are the irreducible multiplets of $G_{SM}$ and the $f_K$ are numerical coefficients. Since we have chosen the DW form for $\langle A \rangle$, Eq.(4) can be written

$$\left( \frac{3}{2} a \left( \frac{P}{M_P} \right) (B - L) \right)^*_K + Z \right) f_K = 0, \quad (6)$$

for all $K$. Since $\langle CC \rangle \neq 0$, not all the $f_K$ vanish. Suppose $f_J$ does not vanish. Then $Z$ is fixed to be $Z = -\frac{3}{2} a \left( \frac{P}{M_P} \right) (B - L)_J$. This, in turn, implies that
There are, therefore, a discrete number of solutions; in fact, four. One of them is \( Z = -\frac{3}{2} a \left( \frac{P}{M_p} \right) \), with \( \langle C \rangle \) pointing in any direction which has \( B - L = 1 \). There is a two-complex-dimensional space of such directions spanned by the “e+” and “Nc” directions; that is, in terms of \( G_{SM} \) quantum numbers, the (1,1,1) and (1,1,0) directions. But actually these are all gauge-equivalent. This is easily checked directly, but is also clear from the fact, which shall presently be seen, that there are no uneaten goldstone modes in this model. Thus, we can take the VEV of \( C \) to lie in the \( SU(5) \)-singlet direction without any loss of generality.

Now we will show that \( \langle C' \rangle = \langle \overline{C}' \rangle = 0 \). That \( C' \) and \( \overline{C}' \) have no VEV in the \( SU(5) \)-singlet direction follows from \( F_Z = F_{\overline{Z}} = 0 \). From the \( SU(5) \)-singlet component of the \( F_C \) and \( F_{\overline{C}} \) equations one has that \( X = 0 \). And from the \( SU(5) \)-non-singlet components of the same equations it follows that \( C' \) and \( \overline{C}' \) have no VEVs in the \( SU(5) \)-non-singlet directions, either. All VEVs have now been fixed except for those of \( PC, P, \overline{P}, \) and \( S \), about which more will be said below.

Knowing the VEVs, one can now read off the Higgsino mass matrices directly from \( W \). For the representations \( K = (3, 2, \frac{1}{6}), (\overline{3}, 1, -\frac{1}{3}), \) and \( (1, 1, +1) \), which are contained in the 10 of \( SU(5) \), one has \( 3 \times 3 \) mass matrices, since such representations exist in the adjoint \( A \) and in the spinors \( C \) and \( C' \). The masses come from the terms in Eq.(3), both through the VEVs of \( A, Z \) and \( \overline{Z} \), and through the VEVs of the spinors \( C \) and \( C' \).

\[
W_{\text{mass,10}}(K) = \begin{pmatrix} A_K & C_K & \overline{C}_K \end{pmatrix} \begin{pmatrix} m_K & 0 & \langle C \rangle \langle P \rangle / \sqrt{2} M_P \\ 0 & 0 & \alpha_K a \langle \overline{P} \rangle / M_P \\ \langle C \rangle \langle P \rangle / \sqrt{2} M_P & \alpha_K a \langle \overline{P} \rangle / M_P & 0 \end{pmatrix} \begin{pmatrix} A_K \\ C_K \\ \overline{C}_K \end{pmatrix}.
\]

(7)

Here \( \alpha_K \equiv \frac{3}{2}((B - L)_K - 1) \), and takes the values \(-1, -2, \) and 0 respectively for \( K = (3, 2, \frac{1}{6}), (\overline{3}, 1, -\frac{1}{3}), \) and \( (1, 1, +1) \). The entry \( m_K \) vanishes for the color-triplet values of \( K \), since \( W_A \) has goldstone modes in those directions, but is non-zero (and in fact equal to \( a^2/2M \)) in the \( (1, 1, +1) \) direction. Thus for each \( K \) the \( 3 \times 3 \) mass matrix has one vanishing eigenvalue, corresponding to a goldstone mode that gets eaten by the Higgs mechanism, and two non-vanishing GUT-scale eigenvalues. One sees also that the massless mode for \( K = (1, 1, +1) \) is purely in the \( C \) direction, as it should be since only
the spinor VEVs break that generator, while for the mass matrices of the color-triplet representations the massless mode is a linear combination of the adjoint and spinor as it should be.

As for the representations \((1, 2, -\frac{1}{2})\) and \((\overline{3}, 1, \frac{1}{2})\) that are contained in the \(\mathbf{5}\) of \(SU(5)\), and their conjugates, they are contained only in the spinors and obtain mass only from the VEVs of \(A\), \(Z\), and \(\overline{Z}\). It is easy to see that the Weak doublets get mass of \(3a \langle P \rangle / M_P\), while the color triplets get mass of \(2a \langle P \rangle / M_P\).

In addition to these, the adjoint contains the \((8, 1, 0)\) and \((1, 3, 0)\), which get mass of \(2a^2 / M\) and \(a^2 / M\) respectively, and the \((3, 2, -\frac{5}{2}) + H.c.\), which get eaten. There are also several singlets of \(G_{SM}\) which get superlarge mass. We have thus seen that no goldstone or pseudo-goldstones are left after symmetry breaking.

From the explicit spectrum given above one can compute the corrections to the low-energy gauge couplings due to the superheavy states. Since \(\sin^2 \theta_W\) and \(\alpha\) are better known, it is now usual to use them as inputs for a prediction of \(\alpha_s(M_Z)\). The minimal \(SU(5)\) SUSY-GUT predicts\(^1\) \(\alpha_s(M_Z) = 0.127 \pm 0.005 \pm 0.002\), where the first error is the uncertainty in the low-energy sparticle spectrum, and the second is the uncertainty in the masses of the top quark and Higgs bosons. A global fit\(^14\) to \(\alpha_s\) from measurements at all energies gives \(\alpha_s(M_Z) = 0.117 \pm 0.005\), while a global fit to all electroweak data\(^15\) gives \(\alpha_s(M_Z) = 0.127 \pm 0.005 \pm 0.002 \pm 0.001\).

Using the notation of Ref. 8, we define \(\epsilon_3 \equiv (\alpha_3(M_G) - \tilde{\alpha}_G) / \tilde{\alpha}_G\), where \(\tilde{\alpha}_G \equiv \alpha_1(M_G) = \alpha_2(M_G)\). (\(M_G\) is here defined to be the scale at which \(\alpha_1\) and \(\alpha_2\) are equal.) To obtain \(\alpha_s(M_Z) \simeq 0.12\) requires, in general, that \(\epsilon_3 \simeq -(2 \text{ to } 3)\%\). In the minimal \(SO(10)\) scheme presented here one finds that \(\epsilon_3 \simeq \frac{3}{5\pi} \tilde{\alpha}_G \ln \left[ \frac{32 \sqrt{2} M_T}{M_G} \right]\), where \(M_T = a^2 / 2 \langle S \rangle\) is the effective color-triplet Higgsino mass that comes into the Higgsino-mediated proton-decay amplitude. \(\epsilon_3\) thus comes out to be \(+0.03\) if \(\tilde{M}_T \simeq 10 M_G\), and \(+0.06\) if \(\tilde{M}_T \simeq 10^3 M_G\) (as is typically necessary if \(\tan \beta\) is large). This is somewhat large, but not necessarily unacceptable, as there can be contributions from other sectors of the theory, such as the quark and lepton sector, that could have the opposite sign.

The stability of the gauge hierarchy requires that certain operators allowed by \(SO(10)\) not be present in \(W\) to sufficiently high order in \(M_P\). These include \(T_1^2, T_1 T_2, \overline{CAC}, \overline{CC}\), or these factors times \(SO(10)\)-singlet
products of fields with non-vanishing VEVs. The first of these types of term can be prevented by a symmetry under which $T_1 \rightarrow e^{i\alpha}T_1$, $T_2 \rightarrow e^{-i\alpha}T_2$, $A \rightarrow A$, and $S \rightarrow e^{2i\alpha}S$. Holomorphy forbids the operator $T_1^2 + \dagger$, of course, but it is necessary also that there be no chiral product of superfields with the same quantum numbers as $S$ and having VEV which is greater than order $M_G/\sqrt{m}$. For example, the VEV of $S$ may not arise from terms like $Y_S(S\bar{S} - m^2)$, as then $T_1^2 S$ would be allowed by symmetry. This is not a problem, however, as $S$ may acquire its VEV radiatively, and there are also ways that it may acquire a tree-level VEV without giving rise to dangerous operators.\footnote{1}

The term $T_1 A T_2$ is an essential ingredient of the DW mechanism, but the term $T_1 T_2 tr A^2$ must be forbidden. This can be done by a $Z_2$ under which $T_1$, $A$, $P$, and $\overline{P}$ are odd and all other fields even.

In fact, all of the destabilizing operators can be forbidden by a $U(1) \times Z_2$ symmetry (or a suitable discrete subgroup of it) under which the fields have the following charges: $A(0, -)$, $P_A(0, +)$, $X(x, +)$, $C(c - \frac{i}{2}, +)$, $\overline{C}(-c - \frac{i}{2}, +)$, $P_C(-\frac{c}{2}, +)$, $P(p, -)$, $\overline{P}(\overline{p}, -)$, $Z(p, +)$, $\overline{Z}(\overline{p}, +)$, $C'(c - p + \frac{i}{2}, +)$, $\overline{C'}(-c - p + \frac{i}{2}, +)$, $T_1(-t, -)$, $T_2(t, +)$, $S(2t, +)$. (The full symmetry of the Higgs superpotential presented above is clearly $U(1)^5 \times Z_2$.) One can now see why the factors $P$ and $\overline{P}$ must appear in Eq. (3). Otherwise, $Z$ and $A$ would have the same quantum numbers and $T_1 Z T_2$ would be allowed.

Destabilizing operators can arise if new fields are introduced to allow some of the singlets to acquire VEVs at tree level. For example, if $\langle P \rangle$ arises from a term $Y(P \bar{P} - m^2)$, then the term $T_1 T_2 Z \bar{P} / M_P$ is allowed (since $\bar{P}$ has the quantum numbers $(-p, -)$). However, as $\langle Z \rangle \sim M_G^3 / M_P$, this gives effectively $(M_G^3 / M_P^2) T_1 T_2$, which in turn gives a “see-saw” contribution to the $\mu$ parameter of order $(M_G^3 / M_P^2)^2 / \langle S \rangle \sim M_G^3 / M_P^3$. This is not only acceptable, but even desirable as a solution to the $\mu$-problem. (Similarly, $\overline{C} A C Z \bar{P} / M_P^2$ would be allowed. It is easy to see that then $\langle A \rangle = \text{diag}(b, b, a, a, a) \otimes i \tau_2$, where $b = O(M_G^3 / M_P^2)$. Again, this gives a “see-saw” contribution to $\mu$ of order $M_G^5 / M_P^4$.)

The VEV of $P_C$ cannot arise, however, from the analogous term $Y_C(P_C \bar{P}_C - M_C^2)$, as that would imply that $\overline{C} C \bar{P}_C^2 / M_P$ was allowed. But a term like $Y_C(P_C^3 \bar{P}_C - M_C^2) / M_P^2$ would be safe. $(\overline{C} C)^3 \bar{P}_C^2 / M_P^3$ would then be allowed, but is harmless.

In conclusion, we have demonstrated a simple supersymmetric SO(10)
model which both breaks SO(10) to the Standard Model and solves the doublet/triplet splitting problem. In one version of the model a $\mu$ term is generated naturally. There is yet one caveat of our approach. In the simplest version there are several flat directions at tree level in which some of the required VEVs must be generated radiatively once SUSY breaking and perhaps supergravity is included.

References

1. S. Dimopoulos, S. Raby, and F. Wilczek, Phys. Rev. D24, 1681 (1981); U. Amaldi, W. deBoer, and H. Furstenau, Phys. Lett. 260B, 447 (1991); P. Langacker and M.-X. Luo, Phys. Rev. D44, 817 (1991); J. Ellis, S. Kelley, and D.V. Nanopoulos, Phys. Lett. 260B, 131 (1991); P. Langacker and N. Polonsky, Phys. Rev. D47, 4028 (1993).

2. J.A. Harvey and M.S. Turner, Phys. Rev. D42, 3344 (1990); G. Lazarides and Q. Shafi, Phys. Lett. B258, 305 (1991); J.A. Harvey and E.W. Kolb, Phys. Rev. D24, 2090 (1981).

3. S.M. Barr, Phys. Rev. D24, 1895 (1981); Phys. Rev. Lett. 64, 353 (1990); K.S. Babu and S.M. Barr, Phys. Rev. Lett. 75, 2088 (1995); G. Anderson, S. Raby, S. Dimopoulos, L.J. Hall, and G.D. Starkman, Phys. Rev. D49, 3660 (1994); L.J. Hall and S. Raby, Phys. Rev. D51, 6524 (1995); V. Lucas and S. Raby, Phys. Rev. D54, 2261 (1996); ibid. D55, 6986 (1997).

4. E. Gildener and S. Weinberg, Phys. Rev. D13, 3333 (1976); E. Gildener, Phys. Rev. D14, 1667 (1976).

5. L. Maiani, in Comptes Rendus de l’Ecole d’Eté de Physiques des Particules, Gif-sur-Yvette, 1979, IN2P3, Paris, 1980, p.3; S. Dimopoulos and H. Georgi, Nucl. Phys B150, 193 (1981); M. Sakai, Z. Phys C11, 153 (1981); E. Witten, Nucl. Phys. B188, 573 (1981).

6. S. Dimopoulos and F. Wilczek, Preprint NSF-ITP-82-07 (1982).
7. K.S. Babu and S.M. Barr, Phys. Rev. D48, 5354 (1993); ibid. D50, 3529 (1994).

8. V. Lucas and S. Raby, ref. 3.

9. K.R. Dienes, Nucl. Phys. B488, 141 (1997).

10. K.S. Babu and S.M. Barr, Phys. Rev. D51, 2463 (1995).

11. J. Hisano, H. Murayama, and T. Yanagida, Phys. Rev. D49, 4966 (1994).

12. V. Lucas and S. Raby, ref. 3; Z. Berezhiani and Z. Tavartkiladze, hep-ph/9612232.

13. S. Urano and R. Arnowitt, hep-ph/9611389.

14. B.R. Webber, in Proceedings 27th International Conference on High Energy Physics, Glasgow, Scotland, 1994, edited by P.J. Bussey and I.G. Knowles (IOP, London, 1995).

15. J. Erler and P. Langacker, Phys. Rev. D52, 441 (1995).