The physics of heavy Z' gauge bosons

Paul Langacker*

School of Natural Sciences, Institute for Advanced Study, Princeton, New Jersey 08540, USA

(Published 7 August 2009)

The U(1)′ symmetry associated with a possible heavy Z' would have profound implications for particle physics and cosmology. The motivations for such particles in various extensions of the standard model, possible ranges for their masses and couplings, and classes of anomaly-free models are discussed. Present limits from electroweak and collider experiments are briefly surveyed, as are prospects for discovery and diagnostic study at future colliders. Implications of a Z' are discussed, including an extended Higgs sector, extended neutralino sector, and solution to the μ problem in supersymmetry; exotic fermions needed for anomaly cancellation; possible flavor changing neutral current effects; neutrino mass; possible Z' mediation of supersymmetry breaking; and cosmological implications for cold dark matter and electroweak baryogenesis.

DOI: 10.1103/RevModPhys.81.1199 PACS number(s): 12.60.Cn, 12.60.Fr, 14.70.Pw

CONTENTS

I. INTRODUCTION

Additional U(1)′ gauge symmetries and associated Z' gauge bosons are one of the best motivated extensions of the standard model (SM). It is not so much that they solve any problems as the fact that it is more difficult to reduce the rank of an extended gauge group containing the standard model than it is to break the non-Abelian factors. As a toy example, consider the gauge group $G=SU(N)$, with $N-1$ diagonal generators. G can be broken by the vacuum expectation value (VEV) of a real

*pgl@ias.edu

adjoint Higgs representation Φ, which can be represented by a Hermitian traceless $N\times N$ matrix

$$\Phi = \sum_{i=1}^{N^2-1} \varphi^i L_i, \quad (1)$$

where φ^i are the real components of Φ and L_i are the fundamental $(N\times N)$ representation matrices. When Φ acquires a VEV $\langle\Phi\rangle$, SU(N) is broken to a subgroup associated with those generators which commute with $\langle\Phi\rangle$. Without loss of generality, $\langle\Phi\rangle$ can be diagonalized by an SU(N) transformation so that the $N-1$ diagonal generators remain unbroken. In special cases some of these may be embedded in unbroken SU(K) subgroups (when K diagonal elements are equal), but the unbroken subgroup always contains at least $U(1)^{N-1}$.

Soon after the proposal of the electroweak $SU(2)\times U(1)_Y$ model there were many suggestions for extended or alternative electroweak gauge theories, some of which involved additional $U(1)'$ factors.[1]

An especially compelling motivation came from the development of grand unified theories larger than the original SU(5) model (Georgi and Glashow 1974), such as those based on SO(10) or E_6 [see, e.g., Robinett and Rosner (1982a, 1982b) and Langacker *et al.* (1984)]. For reviews, see Langacker (1981) and Hewett and Rizzo (1989). These had rank larger than four and could break to $G_{SM}\times U(1)'^n$, $n\geqslant 1$, where $G_{SM}=SU(3)\times SU(2)\times U(1)_Y$ is the standard model gauge group. However, in the original (nonsupersymmetric) versions there was no particular reason for the additional Z' masses to be at the electroweak or TeV scale where they could be directly observed. Similarly, superstring constructions often involve large gauge symmetries which break to $G_{SM}\times U(1)'^n$ in the effective four-dimensional theory (Cvetic and Langacker, 1996a), where some of the $U(1)'$ are nonanomalous. In string theories and in supersymmetric versions of grand unification with extra $U(1)'$'s below the string or GUT scale, both the $U(1)'$ and the $SU(2)\times U(1)_Y$ breaking scales are generally tied to the soft supersymmetry breaking scale (Cvetic and Langacker, 1996a, 1996b, 1997). Therefore, if supersymmetry is observed at the CERN Large Hadron Collider there is a strong motivation that a string or grand united theory (GUT) induced Z' would also have a mass at an observable scale. [An exception to this is when the $U(1)'$ breaking occurs along a flat direction.]

In recent years many TeV scale extensions to the SM have been proposed in addition to supersymmetry, often with the motivation of resolving the fine tuning associated with the quadratic divergence in the Higgs mass. These include various forms of dynamical symmetry breaking (Chivukula and Simmons, 2002; Hill and Simmons, 2003; Chivukula *et al.*, 2004) and little Higgs models (Arkani-Hamed, Cohen and Georgi, 2001; Han *et al.*, 2003, 2006; Perelstein, 2007), which typically involve extended gauge structures, often including new Z' gauge bosons at the TeV scale. Some versions of theories with large extra dimensions allow the standard model gauge bosons to propagate freely in the extra dimensions, implying Kaluza-Klein excitations [see, e.g., Antoniadis (1990), Casalbuoni *et al.* (1999), Masip and Pomarol (1999), Delgado *et al.* (2000), Appelquist *et al.* (2001), Cheng *et al.* (2002), Cheung and Landsberg (2002), Appelquist and Yee (2003), Barbieri *et al.* (2004), and Gogoladze and Macesanu (2006)] of the Z and other standard model gauge bosons, with effective masses of order $R^{-1}\sim 2\ \text{TeV}\times(10^{-17}\ \text{cm}/R)$, where R is the scale of the extra dimension. Such excitations can also occur in Randall-Sundrum models (Randall and Sundrum, 1999) [see, e.g., Hewett *et al.* (2002), Agashe *et al.* (2003, 2007), and Carena, Delgado, *et al.* (2003)].

Other motivations for new Z' bosons, for example, associated with (approximately) hidden sectors of nature, are detailed in Secs. III and V. Extensions of the SM may also involve new TeV scale charged W bosons [see, e.g., Rizzo (2007)], which could couple to either left- or right-handed currents, but the focus of this review will be on Z''s.

The experimental discovery of a new Z' would be exciting, but the implications would be much greater than just the existence of a new vector boson. Breaking the $U(1)'$ symmetry would require an extended Higgs (and neutralino) sector, with significant consequences for collider physics and cosmology (direct searches, the μ problem, dark matter, and electroweak baryogenesis). Anomaly cancellation usually requires the existence of new exotic particles that are vectorlike with respect to the standard model but chiral under $U(1)'$, with several possibilities for their decay characteristics. The expanded Higgs and exotic sectors can modify or maintain the approximate gauge coupling unification of the minimal supersymmetric standard model (MSSM). In some constructions (especially string derived) the $U(1)'$ charges are family nonuniversal, which can lead to flavor changing neutral current (FCNC) effects, for example, in rare B decays. Finally, the decays of a heavy Z' may be a useful production mechanism for exotics and superpartners. The constraints from the $U(1)'$ symmetry can significantly alter the theoretical possibilities for neutrino mass. Finally, $U(1)'$ interactions can couple to a hidden sector, possibly playing a role in supersymmetry breaking or mediation.

Section II of this review discusses basic issues, such as the Z' interactions and properties, $U(1)'$ breaking, anomalies, and ordinary and kinetic mixing between Z

[1]Some examples include Fayet (1977, 1980); Barger and Phillips (1978); Georgi and Weinberg (1978); Mohapatra and Sidhu (1978); Davidson (1979); Barger *et al.* (1980, 1982a, 1982b); Barr and Zee (1980); de Groot *et al.* (1980); Deshpande and Iskandar (1980); Masiero (1980); Rizzo (1980); Zee and Kim (1980); del Aguila and Mendez (1981); Rizzo and Senjanovic (1981); Li and Marshak (1982); Barr (1983). More complete lists of early references can be found in Robinett and Rosner (1982a, 1982b); Langacker *et al.* (1984); Hewett and Rizzo (1989). Previous reviews include Hewett and Rizzo (1989); del Aguila (1994); Cvetic and Godfrey (1995); Cvetic and Langacker (1997); Leike (1999); Rizzo (2006).

and Z'. Section III surveys the large range of models that have been proposed, including the U(1)′-breaking scale; GUT-inspired models; sets of exotics and charges constructed to avoid anomalies; and more exotic possibilities such as ultraweak coupling, low mass, hidden sector, lepto-phobic, intermediate scale, sequential, family nonuniversal, and anomalous U(1)′ models. Section IV outlines the existing constraints from precision electroweak and direct collider searches, as well as prospects for detection and diagnostics of couplings at future colliders. Finally, Sec. V is a survey of the theoretical, collider, and cosmological implications of a possible Z'.

II. BASIC ISSUES

A. Z' couplings

In the standard model the neutral current interactions of the fermions are described by the Lagrangian[2]

$$-L_{NC}^{SM} = gJ_3^\mu W_{3\mu} + g' J_Y^\mu B_\mu = eJ_{em}^\mu A_\mu + g_1 J_1^\mu Z_{1\mu}^0, \qquad (2)$$

where g and g' are the SU(2) and $U(1)_Y$ gauge couplings, $W_{3\mu}$ is the (weak eigenstate) gauge boson associated with the third (diagonal) component of SU(2), and B_μ is the $U(1)_Y$ boson. The currents in the first form are

$$J_3^\mu = \sum_i \bar{f}_i \gamma^\mu [t_{3i_L} P_L + t_{3i_R} P_R] f_i,$$

$$J_Y^\mu = \sum_i \bar{f}_i \gamma^\mu [y_{i_L} P_L + y_{i_R} P_R] f_i, \qquad (3)$$

where f_i is the field of the ith fermion and $P_{L,R} \equiv (1 \mp \gamma^5)/2$ are the left and right chiral projections. t_{3i_L} (t_{3i_R}) is the third component of weak isospin for the left (right) chiral component of f_i. For the known fermions, $t_{3u_L} = t_{3\nu_L} = +\frac{1}{2}$, $t_{3d_L} = t_{3e_L^-} = -\frac{1}{2}$, and $t_{3i_R} = 0$. The weak hypercharges $y_{i_{L,R}}$ are chosen to yield the correct electric charges,

$$t_{3i_L} + y_{i_L} = t_{3i_R} + y_{i_R} = q_i, \qquad (4)$$

where q_i is the electric charge of f_i in units of the positron charge $e > 0$.

Anticipating the spontaneous breaking of SU(2) $\times U(1)_Y$ to the electromagnetic subgroup $U(1)_{em}$ (Sec. II.B), the mass eigenstate neutral gauge bosons in Eq. (2) are the (massless) photon field A_μ and the (massive) $Z_{1\mu}^0 \equiv Z_\mu$, where

$$A_\mu = \sin\theta_W W_{3\mu} + \cos\theta_W B_\mu,$$

$$Z_\mu = \cos\theta_W W_{3\mu} - \sin\theta_W B_\mu, \qquad (5)$$

and the weak angle is $\theta_W \equiv \tan^{-1}(g'/g)$. The new gauge couplings are $e = g\sin\theta_W$ and

$$g_1^2 \equiv g^2 + g'^2 = g^2/\cos^2\theta_W. \qquad (6)$$

The currents in the new basis are

$$J_{em}^\mu = \sum_i q_i \bar{f}_i \gamma^\mu f_i,$$

$$J_1^\mu = \sum_i \bar{f}_i \gamma^\mu [\epsilon_L^1(i) P_L + \epsilon_R^1(i) P_R] f_i, \qquad (7)$$

with the chiral couplings

$$\epsilon_L^1(i) = t_{3i_L} - \sin^2\theta_W q_i, \quad \epsilon_R^1(i) = t_{3i_R} - \sin^2\theta_W q_i. \qquad (8)$$

In the extension to $SU(2)\times U(1)_Y \times U(1)'^n$, $n \geqslant 1$, L_{NC} becomes

$$-L_{NC} = eJ_{em}^\mu A_\mu + \sum_{\alpha=1}^{n+1} g_\alpha J_\alpha^\mu Z_{\alpha\mu}^0, \qquad (9)$$

where g_1, $Z_{1\mu}^0$, and J_1^μ are the gauge coupling, boson, and current of the standard model, respectively. Similarly, g_α and $Z_{\alpha\mu}^0$, $\alpha = 2, \ldots, n+1$, are the gauge couplings and bosons for the additional U(1)′'s. The currents in Eq. (9) are

$$J_\alpha^\mu = \sum_i \bar{f}_i \gamma^\mu [\epsilon_L^\alpha(i) P_L + \epsilon_R^\alpha(i) P_R] f_i$$

$$= \frac{1}{2}\sum_i \bar{f}_i \gamma^\mu [g_V^\alpha(i) - g_A^\alpha(i)\gamma^5] f_i. \qquad (10)$$

The chiral couplings $\epsilon_{L,R}^\alpha(i)$, which may be unequal for a chiral gauge symmetry, are the $U(1)_\alpha$ charges of the left- and right-handed components of fermion f_i, and $g_{V,A}^\alpha(i) = \epsilon_L^\alpha(i) \pm \epsilon_R^\alpha(i)$ are the corresponding vector and axial couplings, respectively.

Frequently, it is more convenient to instead specify the $U(1)_\alpha$ charges of the left chiral components of both the fermion f and the antifermion (conjugate) f^c, denoted $Q_{\alpha f}$ and $Q_{\alpha f^c}$, respectively. The two sets of charges are simply related,

$$\epsilon_L^\alpha(f) = Q_{\alpha f}, \quad \epsilon_R^\alpha(f) = -Q_{\alpha f^c}. \qquad (11)$$

For example, in the SM one has $Q_{1u} = \frac{1}{2} - \frac{2}{3}\sin^2\theta_W$ and $Q_{1u^c} = +\frac{2}{3}\sin^2\theta_W$.

The additional gauge couplings and charges, as well as the gauge boson masses and mixings, are extremely model dependent. The gauge couplings and charges are not independent, i.e., one can always replace g_α by $\lambda_\alpha g_\alpha$ provided the charges Q_α are all simultaneously scaled by $1/\lambda_\alpha$. Usually, the charges are normalized by some convenient convention.

The three- and four-point gauge interactions of a complex SU(2) scalar multiplet ϕ can be read off from the kinetic term $L_\phi^{kin} = (D^\mu\phi)^\dagger D_\mu \phi$. The diagonal (neutral current) part of the gauge covariant derivative of an individual field ϕ_i is

[2]We largely follow the formalism and conventions in Durkin and Langacker (1986) and Langacker and Luo (1992).

$$D_\mu \phi_i = \left(\partial_\mu + ieq_i A_\mu + i \sum_{\alpha=1}^{n+1} g_\alpha Q_{\alpha i} Z^0_{\alpha\mu} \right) \phi_i, \tag{12}$$

where q_i and $Q_{\alpha i}$ are the electric and $U(1)_\alpha$ charges of ϕ_i, respectively. For the SM part, $l_i = 0, \frac{1}{2}, 1, \ldots$ labels the SU(2) representation, l_{3i} is the third component of weak isospin, the weak hypercharge is $y_i = q_i - l_{3i}$, and $Q_{1i} = l_{3i} - \sin^2\theta_W q_i$. Thus, for the neutral component ϕ^0 of the Higgs doublet $\phi = \binom{\phi^+}{\phi^0}$ one has $t_{\phi^0} = -t_{3\phi^0} = y_{\phi^0} = \frac{1}{2}$.

B. Masses and mass mixings

We assume that electrically neutral scalar fields ϕ_i acquire VEVs, so A_μ remains massless, while the $Z^0_{\alpha\mu}$ fields develop a mass term $L_Z^{\text{mass}} = \frac{1}{2} M^2_{\alpha\beta} Z^0_{\alpha\mu} Z^{0\mu}_\beta$, where

$$M^2_{\alpha\beta} = 2 g_\alpha g_\beta \sum_i Q_{\alpha i} Q_{\beta i} |\langle \phi_i \rangle|^2. \tag{13}$$

$M^2_{11} \equiv M^2_{Z^0}$ would be the (tree-level) Z mass in the SM limit in which the other Z^0's and their mixing can be ignored. If the only Higgs fields are SU(2) doublets (or singlets), as in the SM or the MSSM, then

$$M^2_{Z^0} = \tfrac{1}{2} g_1^2 \sum_i |\langle \phi_i \rangle|^2 = \tfrac{1}{4} g_1^2 \nu^2 = \frac{M_W^2}{\cos^2\theta_W}, \tag{14}$$

where $\nu^2 = 2\sum_i |\langle \phi_i \rangle|^2 \sim (\sqrt{2} G_F)^{-1} \sim (246 \text{ GeV})^2$ is the square of the weak scale and G_F is the Fermi constant. The observed Z mass strongly constrains either higher-dimensional Higgs VEVs or Z-Z' mixing (Yao *et al.*, 2006), but in principle they could compensate and should both be considered. Allowing a general Higgs structure, one has

$$M^2_{Z^0} = \frac{g_1^2}{4\sqrt{2} G_F \rho_0} = \frac{M_W^2}{\rho_0 \cos^2\theta_W}, \tag{15}$$

where

$$\rho_0 \equiv \frac{\sum_i (t_i^2 - t_{3i}^2 + t_i) |\langle \phi_i \rangle|^2}{\sum_i 2 t_{3i}^2 |\langle \phi_i \rangle|^2} \xrightarrow[\text{doublets, singlets}]{} 1. \tag{16}$$

Diagonalizing the mass matrix, Eq. (13), one obtains $n+1$ (usually) massive eigenstates $Z_{\alpha\mu}$ with mass M_α,

$$Z_{\alpha\mu} = \sum_{\beta=1}^{n+1} U_{\alpha\beta} Z^0_{\beta\mu}, \tag{17}$$

where U is an orthogonal mixing matrix. It is straightforward to show that the mass-squared eigenvalues are always non-negative. From Eqs. (9) and (17) $Z_{\alpha\mu}$ couples to $\sum_\beta g_\beta U_{\alpha\beta} J^\mu_\beta$.

The most studied case is $n=1$. Writing $Q_i \equiv Q_{2i}$, the mass matrix is

$$M^2_{Z\text{-}Z'} = \begin{pmatrix} 2g_1^2 \sum_i t_{3i}^2 |\langle \phi_i \rangle|^2 & 2 g_1 g_2 \sum_i t_{3i} Q_i |\langle \phi_i \rangle|^2 \\ 2 g_1 g_2 \sum_i t_{3i} Q_i |\langle \phi_i \rangle|^2 & 2 g_2^2 \sum_i Q_i^2 |\langle \phi_i \rangle|^2 \end{pmatrix}$$
$$\equiv \begin{pmatrix} M^2_{Z^0} & \Delta^2 \\ \Delta^2 & M^2_{Z'} \end{pmatrix}. \tag{18}$$

As an example, many U(1)$'$ models involve an SU(2) singlet S and two Higgs doublets

$$\phi_u = \begin{pmatrix} \phi_u^0 \\ \phi_u^- \end{pmatrix}, \quad \phi_d = \begin{pmatrix} \phi_d^+ \\ \phi_d^0 \end{pmatrix}, \tag{19}$$

with U(1)$'$ charges $Q_{S,u,d}$. Then

$$M^2_{Z^0} = \tfrac{1}{4} g_1^2 (|\nu_u|^2 + |\nu_d|^2),$$
$$\Delta^2 = \tfrac{1}{2} g_1 g_2 (Q_u |\nu_u|^2 - Q_d |\nu_d|^2),$$
$$M^2_{Z'} = g_2^2 (Q_u^2 |\nu_u|^2 + Q_d^2 |\nu_d|^2 + Q_S^2 |s|^2), \tag{20}$$

where $\nu_{u,d} \equiv \sqrt{2} \langle \phi^0_{u,d} \rangle$, $s = \sqrt{2} \langle S \rangle$, and $\nu^2 = (|\nu_u|^2 + |\nu_d|^2) \sim (246 \text{ GeV})^2$.

The eigenvalues of a general $M^2_{Z\text{-}Z'}$ are

$$M^2_{1,2} = \tfrac{1}{2} [M^2_{Z^0} + M^2_{Z'} \mp \sqrt{(M^2_{Z^0} - M^2_{Z'})^2 + 4\Delta^4}], \tag{21}$$

and U is the rotation

$$U = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}, \tag{22}$$

with

$$\theta = \frac{1}{2} \arctan\left(\frac{2\Delta^2}{M^2_{Z^0} - M^2_{Z'}} \right). \tag{23}$$

θ is related to the masses by

$$\tan^2\theta = \frac{M^2_{Z^0} - M_1^2}{M_2^2 - M^2_{Z^0}}. \tag{24}$$

An important limit is $M_{Z'} \gg (M_{Z^0}, |\Delta|)$, which typically occurs because an SU(2) singlet field (such as S in the example) has a large VEV and contributes only to $M_{Z'}$. One then has

$$M_1^2 \sim M^2_{Z^0} - \frac{\Delta^4}{M^2_{Z'}} \ll M_2^2, \quad M_2^2 \sim M^2_{Z'} \tag{25}$$

and

$$\theta \sim -\frac{\Delta^2}{M^2_{Z'}} \sim C \frac{g_2}{g_1} \frac{M_1^2}{M_2^2} \quad \text{with } C = -\frac{\sum_i t_{3i} Q_i |\langle \phi_i \rangle|^2}{\sum_i t_{3i}^2 |\langle \phi_i \rangle|^2}. \tag{26}$$

C is model dependent, but typically $|C| \lesssim O(1)$. From Eqs. (24)–(26) one sees that both $|\theta|$ and the downward shift $(M_{Z^0} - M_1)/M_{Z^0}$ are of order M_1^2/M_2^2. Generalizations of these results for $n>1$ extra U(1)s are given in Langacker (1984).

C. Anomalies and exotics

A symmetry is chiral if it acts differently on the left- and right-handed fermions, and nonchiral (or vector) otherwise. Thus, a chiral $U(1)'$ has $Q_{\alpha f} \neq -Q_{\alpha f^c}$ for at least one f, which is also referred to as chiral. Even for a chiral symmetry, some of the fermions may be nonchiral. If a given fermion pair is nonchiral with respect to all symmetries then an elementary mass term $-L_m = m_f \bar{f}_L f_R + \text{H.c.}$ is allowed, where m_f could be arbitrarily large. Such a term is forbidden for a chiral fermion, whose mass is only generated when the symmetry is broken. For example, if the symmetries allow the Yukawa coupling

$$-L_{\text{Yuk}} = \lambda_f \varphi \bar{f}_L f_R + \text{H.c.}, \tag{27}$$

where φ is charged under the symmetry, then an effective mass $\lambda_f \langle \varphi \rangle$ is generated when φ acquires a VEV. Assuming $\lambda_f \lesssim 1$, m_f cannot be larger than the symmetry breaking scale $\langle \varphi \rangle$. In the SM, the ordinary quarks and leptons are chiral under both SU(2) and $U(1)_Y$, and φ is the Higgs doublet. Similar constraints apply to new fields occurring in $U(1)'$ models, which are frequently *quasichiral*, i.e., nonchiral under the SM but chiral under $U(1)'$.

Consistency of a low-energy gauge theory requires the absence of triangle anomalies, including mixed gauge-gravitational ones.[3] For the SM, the nontrivial conditions are

$$\sum_f Y_f = 0, \quad \sum_f Y_f^3 = 0,$$

$$\sum_{f \in 3,3^*} Y_f = 0, \quad \sum_{f \in 2} Y_f = 0, \tag{28}$$

where the sum extends over all left-handed fermions ($Y_f = y_{f_L}$) and antifermions ($Y_{f^c} = -y_{f_R}$). The first condition is the mixed anomaly; the sum is over color triplets and antitriplets in the third $[SU(3)^2 Y]$ condition; and the sum is over SU(2) doublets in the last $[SU(2)^2 Y]$ condition. The sum includes counting factors of 3 for families, 3 for color triplets, and 2 for SU(2) doublets since SU(3) and SU(2) commute with hypercharge and additional $U(1)'$'s. For example, the second $[Y^3]$ condition is

$$3[6Y_Q^3 + 2Y_L^3 + 3Y_{u^c}^3 + 3Y_{d^c}^3 + Y_{e^c}^3] = 0, \tag{29}$$

where Q and L refer to quark and lepton doublets, respectively. This is satisfied by a cancellation between quark and lepton terms. One also requires the absence of an $[SU(3)^3]$ anomaly. In the SM this is achieved automatically because there are equal numbers of quarks and antiquarks. With an additional $U(1)'$ with charge Q_2, there are additional conditions obtained from Eq. (28) by replacing Y by Q_2. There are also mixed $[YU(1)']$ conditions $\Sigma_f Y Q_2^2 = \Sigma_f Y^2 Q_2 = 0$. For $n>1$ additional $U(1)'$'s there are similar conditions for every $Q_\beta, \beta \geq 2$, as well as $\Sigma_f Y Q_\alpha Q_\beta = \Sigma_f Q_\alpha Q_\beta Q_\gamma = 0$. All of these sums include any extra chiral fermions in the theory, such as the superpartners of Higgs scalars in supersymmetry. Nonchiral fermion pairs cancel.

Even for a single $U(1)'$ it is easy to show that the anomaly conditions cannot be satisfied by the SM fermions alone if the $U(1)'$ charges are the same for all three families, except for the trivial case $Q_2 = 0$. ($Q_2 = cY$ is also possible, but this is equivalent to $Q_2 = 0$ after performing a rotation on B_μ and $Z_{2\mu}^0$.) Thus, almost all $U(1)'$ constructions involve additional fermions, known as exotics. These may be singlets under the SM gauge group, such as a singlet right-handed neutrino, or they may carry nontrivial SM quantum numbers. Precision electroweak constraints strongly restrict, but do not entirely exclude, the possibilities of new fermions chiral under $SU(2) \times U(1)_Y$ (Yao *et al.*, 2006), so such exotics are usually assumed to be quasichiral, e.g., both left- and right-handed components might be SU(2) doublets, or both might be singlets. A typical example is a new SU(2)-singlet heavy down-type quark D with $q_D = -1/3$ and its partner D^c.

One can introduce vector pairs that are charged but nonchiral under both the SM and $U(1)'$. These do not contribute to the anomaly conditions, but contribute to the renormalization group equations (RGE) for the gauge couplings, and may also be relevant to the decays of exotics.

If two $U(1)'$ charges $Q_{\alpha,\beta}$ (one of these can be Y) are both generators of a simple underlying group, then one expects them to be orthogonal, i.e., $\Sigma_f Q_{\alpha f} Q_{\beta f} = 0$ for $\alpha \neq \beta$, with a corresponding condition for the scalar charges. However, this condition need not hold without such an embedding or for a more complicated one, or it could be violated due to kinetic mixing (Sec. II.D). Furthermore, all fermions, including the nonchiral ones, contribute to the orthogonality condition. In particular, an apparent violation of orthogonality could be due to the fact that the contributions of a very heavy vector pair (or of heavy scalars) have not been taken into account. There is always some freedom to perform rotations on the gauge fields $Z_{\alpha\mu}^0$, e.g., to make the $U(1)'$ charges orthogonal (at least with respect to the fermions, in a nonsupersymmetric theory). However, such a rotated basis may not coincide with either the mass or kinetic eigenstates.

D. Kinetic mixing

The most general kinetic energy term for the two gauge bosons $Z_{\alpha\mu}^0$ and $Z_{\beta\mu}^0$ in $U(1)_\alpha \times U(1)_\beta$ is

[3]See Sec. III for the role of anomalous $U(1)'$'s.

$$L_{\rm kin} \to -\frac{c_\alpha}{4}F^{0\mu\nu}_\alpha F^0_{\alpha\mu\nu} - \frac{c_\beta}{4}F^{0\mu\nu}_\beta F^0_{\beta\mu\nu} - \frac{c_{\alpha\beta}}{2}F^{0\mu\nu}_\alpha F^0_{\beta\mu\nu}, \tag{30}$$

where $F^0_{\alpha\mu\nu}=\partial_\mu Z^0_{\alpha\nu}-\partial_\nu Z^0_{\alpha\mu}$. One can put the first two terms into canonical form $c_\alpha=c_\beta=1$ by rescaling the fields, and take $c_{\alpha\beta}=\sin\chi$. Since U(1) field strengths are invariant, the cross (kinetic mixing) term does not spoil the gauge invariance (Holdom, 1986). Even if $\chi=0$ at tree level, it can be generated by loop effects if there are particles in the theory that are simultaneously charged under both U(1)'s (Holdom, 1986; Matsuoka and Suematsu, 1986; del Aguila *et al.*, 1988; Foot and He, 1991; del Aguila, Masip, and Perez-Victoria, 1995; Babu *et al.*, 1996, 1998). The mixing term can be cast as a cross term in the RGE for the gauge couplings, with a coefficient proportional to $\Sigma_{m_f<\mu}Q_{\alpha f}Q_{\beta f}$, where μ is the RGE scale, with corresponding contributions from scalars. Even for orthogonal charges the sum at lower mass scales may be nonzero due to the decoupling of heavy particles. Such RGE effects are usually of order a few percent in χ, but could be larger if there are many decoupled states (Babu *et al.*, 1996). A nonzero χ can also be generated by string loop effects in superstring theory. These contributions are small in the heterotic constructions considered in Dienes *et al.* (1997). However, if one of the U(1) factors is broken in a hidden sector at a large scale, the associated D terms could propagate this scale to the ordinary sector by kinetic mixing, destabilizing the supersymmetry breaking scale, and leading to negative mass-square scalars (Dienes *et al.*, 1997).

Now, consider the consequences of kinetic mixing for a single extra U(1)$'$, i.e., $\alpha=1$, $\beta=2$. L_{kin} can be put in canonical form (for $c_{1,2}=1$, $c_{12}=\sin\chi$) by defining

$$\begin{pmatrix} Z^0_{1\mu} \\ Z^0_{2\mu}\end{pmatrix} = \begin{pmatrix} 1 & -\tan\chi \\ 0 & 1/\cos\chi\end{pmatrix}\begin{pmatrix} \hat{Z}^0_{1\mu} \\ \hat{Z}^0_{2\mu}\end{pmatrix} \equiv V\begin{pmatrix} \hat{Z}^0_{1\mu} \\ \hat{Z}^0_{2\mu}\end{pmatrix}, \tag{31}$$

where V is nonunitary. In the new $\hat{Z}$ basis, the mass matrix in Eq. (18) becomes $V^T M^2_{Z\text{-}Z'}V$, which can be diagonalized by an orthogonal matrix U^T. Similarly, the interaction term in Eq. (9) becomes

$$\begin{aligned}(g_1 J_1^\mu g_2 J_2^\mu)\begin{pmatrix} Z^0_{1\mu} \\ Z^0_{2\mu}\end{pmatrix} &\equiv \mathcal{J}^T\begin{pmatrix} Z^0_{1\mu} \\ Z^0_{2\mu}\end{pmatrix} \to \mathcal{J}^T V\begin{pmatrix} \hat{Z}^0_{1\mu} \\ \hat{Z}^0_{2\mu}\end{pmatrix} \\ &= \mathcal{J}^T V U^T \begin{pmatrix} Z_{1\mu} \\ Z_{2\mu}\end{pmatrix},\end{aligned} \tag{32}$$

where $Z_{1,2}$ are the mass eigenstates. These transformations have been analyzed in detail by Babu *et al.* (1998). The essential feature can be seen for $\Delta^2=0$ in Eq. (18), for which

$$V^T M^2_{Z\text{-}Z'}V = \begin{pmatrix} M^2_{Z^0} & -M^2_{Z^0}\tan\chi \\ -M^2_{Z^0}\tan\chi & M^2_{Z^0}\tan^2\chi + M^2_{Z'}/\cos^2\chi\end{pmatrix}. \tag{33}$$

One sees immediately that for $M^2_{Z^0}=0$ there is a zero eigenvalue, even for large χ, i.e., any shift in the lighter mass induced by kinetic mixing is proportional to the light mass and therefore small. In fact, for $|M^2_{Z^0}|\ll|M^2_{Z'}|$, $\Delta^2=0$, and $|\chi|\ll 1$ one has $M_1^2\sim M^2_{Z^0}-M^4_{Z^0}\chi^2/M^2_{Z'}$, a negligible shift. The only significant effect in this limit is that the couplings become

$$g_1 J_1^\mu Z_{1\mu} + (g_2 J_2^\mu - g_1\chi J_1^\mu)Z_{2\mu}, \tag{34}$$

i.e., the coupling of the heavy boson is shifted to include a small component proportional to J_1. The light boson couplings are not affected to this order. One must still include the further effects of mass mixing ($\Delta^2\neq 0$). For $|\Delta^2|\lesssim|M^2_{Z^0}|\ll|M^2_{Z'}|$ and $|\chi|\ll 1$,

$$M_1^2 \sim M^2_{Z^0} - \frac{(\Delta^2 - M^2_{Z^0}\chi)^2}{M^2_{Z'}} \sim M^2_{Z^0} - \hat{\theta}^2 M^2_{Z'}, \tag{35}$$

where

$$\hat{\theta} \equiv \frac{-\Delta^2 + M^2_{Z^0}\chi}{M^2_{Z'}}. \tag{36}$$

This is of the same form as Eq. (25) except that the effective mixing angle is shifted from Eq. (26) by the kinetic mixing. The interactions are just the rotation by $\hat{\theta}$ of those in Eq. (34).

In a supersymmetric theory the charges in the U(1)$_2$ D terms are also shifted, $g_2Q_2\to g_2Q_2-g_1\chi Q_1$. There can also be kinetic mixing between the U(1)$'$ gauginos (Suematsu, 1999), with consequences analogous to those for the gauge bosons.

E. One and two Higgs doublets, supersymmetry, and the μ problem

1. Higgs doublets

The standard model involves a single Higgs doublet $\phi=\binom{\phi^+}{\phi^0}$, which has Yukawa couplings (ignoring family indices)

$$\begin{aligned}-L_{\rm Yuk} &= h_d\bar{Q}_L\phi d_R + h_u\bar{Q}_L\tilde{\phi}u_R + h_e\bar{L}_L\phi e^-_R \\ &\quad + h_\nu\bar{L}_L\tilde{\phi}\nu_R + \text{H.c.},\end{aligned} \tag{37}$$

where $Q_L\equiv\binom{u_L}{d_L}$, $L_L\equiv\binom{\nu_L}{e^-_L}$, and ν_R is the right-handed (SM-singlet) neutrino. The tilde field is defined by

$$\tilde{\phi} \equiv i\sigma^2\phi^* = \begin{pmatrix}\phi^{0*} \\ -\phi^-\end{pmatrix}, \tag{38}$$

where σ^2 is the second Pauli matrix. It is essentially the Hermitian conjugate of ϕ, but transforms as a **2** rather than a **2*** under SU(2), and has $y=-1/2$. A single doublet suffices for the SM, but in many extensions, including supersymmetry and many U(1)′ models, the $\tilde{\phi}$ couplings are not allowed. One must introduce a second (independent) doublet ϕ_u, as in Eq. (19), which plays the role of $\tilde{\phi}$, while ϕ_d plays that of ϕ.

In supersymmetric models it is convenient to work entirely in terms of (left) chiral superfields, such as Q, L, u^c, d^c, e^+, and the SM singlet ν^c which is conjugate to ν_R (we do not distinguish between chiral superfields and their components in our notation—the context should always make the meaning clear). Furthermore, supersymmetry (and anomaly constraints) require two Higgs doublets $H_u=\binom{H_u^+}{H_u^0}$ and $H_d=\binom{H_d^0}{H_d^-}$ with $Y_{H_{u,d}}=\pm 1/2$, defined so that the MSSM superpotential is

$$W=\mu H_u H_d - h_d Q H_d d^c + h_u Q H_u u^c - h_e L H_d e^+ + h_\nu L H_u \nu^c. \tag{39}$$

Doublets are contracted according to $H_u H_d \equiv \epsilon^{ab} H_{Ua} H_{db}$, etc., where $\epsilon^{12}=-\epsilon^{21}=1$. The two sets of Higgs doublets are related by $H_{u,d}=\mp\tilde{\phi}_{u,d}$. The superpotential in Eq. (39) assumes that R parity is conserved. Some U(1)′ models enforce this automatically, as will be mentioned in Sec. III.B.

2. Nonholomorphic terms

In some U(1)′ extensions of the MSSM, some of the Yukawa couplings in Eq. (39) may be forbidden by the U(1)′ gauge symmetry. In some cases, however, the operators involving the *wrong* Higgs field, such as $Q\tilde{H}_u d^c$ or $L\tilde{H}_u e^+$, may be U(1)′ invariant. Such nonholomorphic operators are not allowed in W by supersymmetry, but could be present in the Kähler potential, where they would lead to corresponding nonholomorphic soft terms (Borzumati *et al.*, 1999) for the scalar squarks and sleptons. These then lead to fermion masses at one loop by gluino or neutralino exchange. However, in most supersymmetry breaking schemes it is difficult to generate a large enough effective Yukawa potential (Martin, 1997) because the nonholomorphic soft terms have an additional suppression (compared to the usual soft SUSY breaking scale of $M_{\rm SUSY}\sim 1$ TeV) of $M_{\rm SUSY}/M_{\rm med}$, where $M_{\rm med}\gg M_{\rm SUSY}$ is the SUSY mediation scale (such as the Planck scale for supergravity mediation).

3. The μ problem

One difficulty with the MSSM is the μ problem (Kim and Nilles, 1984), i.e., the supersymmetric Higgs mass μ in Eq. (39) could be arbitrarily large, but phenomenologically needs to be of the same order as the soft supersymmetry breaking terms. In many supersymmetric U(1)′ models this problem is solved because an elementary μ term is forbidden by the U(1)′, but a trilinear $W_\mu=\lambda_S S H_u H_d$ is allowed, where S is a singlet under the SM but charged under the U(1)′. Then, a dynamical effective $\mu_{\rm eff}=\lambda_S\langle S\rangle$ is generated that is related to the scale of U(1)′ breaking (Suematsu and Yamagishi, 1995; Cvetic and Langacker, 1996a; Cvetic *et al.*, 1997), as discussed in Sec. V.A.1. This mechanism can also be associated with discrete or other symmetries (Accomando *et al.*, 2006). An alternative solution, the Giudice-Masiero mechanism, generates μ through a nonrenormalizable operator in the Kähler potential (Giudice and Masiero, 1988). It is especially useful when an elementary μ term is allowed by the low energy symmetries of the theory, but is forbidden by the underlying string construction. This mechanism can also be used to generate mass for vector pairs in U(1)′ theories.

III. MODELS

There are enormous numbers of U(1)′ models, and it is only possible to touch on the major classes and issues here. The models are distinguished by (a) the coupling constants g_α, which are often assumed to be of electroweak strength, but could be larger or smaller. (b) The U(1)′ breaking scale. In some scenarios this is arbitrary, with no good reason to expect it to be around the TeV scale. However, in supersymmetric models it is usually at the TeV scale, unless the breaking is associated with an F and D flat direction, when it could be much larger. The TeV scale is also expected when the U(1)′ is associated with alternative models of electroweak breaking. String constructions usually imply some Z''s close to the string scale, and often involve lighter ones as well. Finally, a Z' could actually be lighter than the electroweak scale if its couplings to the SM fields are small. (c) Other critical issues are the charges of the SM fermions and Higgs doublet, and whether the fermion charges are family universal; the type of scalar responsible for the U(1)′ breaking; whether additional exotic fields are needed to cancel anomalies; whether the theory is supersymmetric (so that the Higgs superpartners must be included in the anomaly considerations); whether the Yukawa couplings of the ordinary fermions are allowed by the U(1)′ symmetry; and whether other couplings, such as those associated with the supersymmetric μ parameter, R-parity violation, and Majorana neutrino masses are allowed.

A. Canonical examples

1. Sequential model

The *sequential* $Z_{\rm SM}$ boson is defined to have the same couplings to fermions as the SM Z boson. The $Z_{\rm SM}$ is not expected in the context of gauge theories unless it has different couplings to exotic fermions, or if it occurs as an excited state of the ordinary Z in models with extra dimensions at the weak scale. However, it serves as a useful reference case when comparing constraints from various sources.

2. Models based on T_{3R} and B-L

One of the simplest and most common classes of models involves $SU(2)\times U(1)_{3R}\times U(1)_{BL}$, where the $U(1)_{3R}$ generator T_{3R} is $\frac{1}{2}$ for (u_R,ν_R), $-\frac{1}{2}$ for (d_R,e_R^-), and 0 for f_L; and the $U(1)_{BL}$ generator is $T_{BL}\equiv\frac{1}{2}(B\text{-}L)$, where B

TABLE I. Charges of the left-chiral components of the fermions in the models based on T_{3R} and $T_{BL}=(B-L)/2$. The charges are normalized so that $g_2=\sqrt{\frac{5}{3}}g\tan\theta_W$. Q^{LR} is defined in Eq. (42), and $Q^{YBL}\equiv b(zY+T_{BL})$. Q^{LR} is a special case of Q^{YBL} for $b^2z(1+z)=-3/5$. α and (b,z) are free parameters, with $\alpha=1.53$ for left-right symmetry and $\alpha\sim 0.7-0.9$ for most SO(10) models.

	T_{3R}	T_{BL}	Y	$\sqrt{\frac{5}{3}}Q^{LR}$	$\frac{1}{b}Q^{YBL}$
Q	0	$\frac{1}{6}$	$\frac{1}{6}$	$-\frac{1}{6\alpha}$	$\frac{1}{6}(z+1)$
u_L^c	$-\frac{1}{2}$	$-\frac{1}{6}$	$-\frac{2}{3}$	$-\frac{\alpha}{2}+\frac{1}{6\alpha}$	$-\frac{2}{3}z-\frac{1}{6}$
d_L^c	$\frac{1}{2}$	$-\frac{1}{6}$	$\frac{1}{3}$	$\frac{\alpha}{2}+\frac{1}{6\alpha}$	$\frac{1}{3}z-\frac{1}{6}$
L_L	0	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2\alpha}$	$-\frac{1}{2}(z+1)$
e_L^+	$\frac{1}{2}$	$\frac{1}{2}$	1	$\frac{\alpha}{2}-\frac{1}{2\alpha}$	$z+\frac{1}{2}$
ν_L^c	$-\frac{1}{2}$	$\frac{1}{2}$	0	$-\frac{\alpha}{2}-\frac{1}{2\alpha}$	$\frac{1}{2}$

(L) is baryon (lepton) number; and ν_R are right-handed neutrinos (see Table I). T_{3R} and T_{BL} are related to weak hypercharge by $Y=T_{3R}+T_{BL}$. T_{3R} occurs in left-right symmetric models based on the group $G_{LR}\equiv SU(2)_L\times SU(2)_R\times U(1)_{BL}$ [for a review, see Mohapatra (1986)] and in SO(10) models (which contain G_{LR}) (Langacker, 1981; Hewett and Rizzo, 1989). The Higgs doublet ϕ can be assigned $T_{3R}=\frac{1}{2}$ and $T_{BL}=0$. However, in the G_{LR} or SO(10) embeddings (or in supersymmetric versions), there are two Higgs doublets $\phi_{u,d}$, as defined in Eq. (19), with $T_{3R}=\mp\frac{1}{2}$. All of these versions are anomaly-free after including the three ν_R.

For these models, the fermion neutral current couplings are

$$-L_{\rm NC}=gJ_{3L}W_{3L}+g_RJ_{3R}W_{3R}+g_{BL}J_{BL}W_{BL}, \qquad (40)$$

where $J_{3L}\equiv J_3$, J_{3R} and J_{BL} are the currents corresponding to T_{3R} and T_{BL}, the g's and W's are the coupling constants and gauge bosons, and the Lorentz indices have been suppressed.

We anticipate that $U(1)_{3R}\times U(1)_{BL}$ will be broken to $U(1)_Y$ at a scale $M_{Z'}\gg M_{Z^0}$, so it is convenient to first transform the gauge bosons W_{3R} and W_{BL} to a new basis B and Z_2^0, where B is identified with the SM $U(1)_Y$ boson,

$$\begin{aligned}-L_{\rm NC}&=gJ_{3L}W_{3L}+g'J_YB+g_2J_2Z_2^0\\&=eJ_{\rm em}A+g_1J_1Z_1^0+g_2J_2Z_2^0,\end{aligned} \qquad (41)$$

as in Eq. (9). We first assume that the gauge kinetic terms are canonical, i.e., with unit strength and no kinetic mixing, so that orthogonal transformations on the three gauge bosons will leave the kinetic terms invariant. Taking $B\equiv\cos\gamma W_{3R}+\sin\gamma W_{BL}$ and choosing γ so that B couples to $g'Y$, one finds $1/g'^2=1/g_R^2+1/g_{BL}^2$, and that the orthogonal gauge boson $Z_2^0=\sin\gamma W_{3R}-\cos\gamma W_{BL}$ couples to the current J_2 associated with the charge

$$Q^{LR}=\sqrt{\frac{3}{5}}\left[\alpha T_{3R}-\frac{1}{\alpha}T_{BL}\right], \qquad (42)$$

where

$$\alpha=\frac{g_R}{g_{BL}}=\sqrt{\kappa^2\cot^2\theta_W-1}, \qquad (43)$$

with $\kappa\equiv g_R/g$. The coupling has been normalized to

$$g_2=\sqrt{\frac{5}{3}}g\tan\theta_W\sim 0.46 \qquad (44)$$

for later convenience.

One interesting case is when G_{LR} survives down to the TeV scale. This is usually studied assuming a left-right symmetry under the interchange of the two SU(2) factors (Mohapatra, 1986), in which case $g_R=g$ and $\alpha\sim 1.53$ for $\sin^2\theta_W\sim 0.23$. Two forms of the model are often considered. In both cases, the Higgs doublets $\phi_{u,d}$ responsible for fermion mass transform as $(2,2)_0$ under $SU(2)_L\times SU(2)_R$, where the subscript is the T_{BL} charge. In one class, an additional doublet pair $\delta_{R,L}$ transforming as $(1,2)_{1/2}+(2,1)_{1/2}$ is introduced, with the VEV of δ_R^0 breaking G_{LR} to the SM. In the other, one instead introduces a triplet pair $\Delta_{R,L}$ transforming as $(1,3)_1+(3,1)_1$. The Δ_R^0 VEV not only breaks G_{LR} but also leads to a large Majorana mass for the ν_R and therefore a small ν_L mass by the seesaw mechanism (Mohapatra and Senjanovic, 1980). The low-scale left-right model also implies a new $W_R^\pm$ which couples to right-handed

TABLE II. Decomposition of the E_6 fundamental representation of left-handed fermions **27** under SO(10) and SU(5), and their $U(1)_\chi$, $U(1)_\psi$, $U(1)_\eta$, inert $U(1)_I$, neutral-N $U(1)_N$, and secluded sector $U(1)_S$ charges. A general model in this class has charge $Q_2 = \cos\theta_{E_6} Q_\chi + \sin\theta_{E_6} Q_\psi - \epsilon Y$, where ϵ can result from kinetic mixing, and coupling $g_2 = \sqrt{\frac{5}{3}} g \tan\theta_W \lambda_g^{1/2}$, where λ_g is usually of $O(1)$.

SO(10)	SU(5)	$2\sqrt{10}Q_\chi$	$2\sqrt{6}Q_\psi$	$2\sqrt{15}Q_\eta$	$2Q_I$	$2\sqrt{10}Q_N$	$2\sqrt{15}Q_S$
16	$10(u,d,u^c,e^+)$	−1	1	−2	0	1	−1/2
	$5^*(d^c,\nu,e^-)$	3	1	1	−1	2	4
	ν^c	−5	1	−5	1	0	−5
10	$5(D,H_u)$	2	−2	4	0	−2	1
	$5^*(D^c,H_d)$	−2	−2	1	1	−3	−7/2
1	$1S$	0	4	−5	−1	5	5/2

currents and can mix with the SM $W^\pm$. Strategies for determining the symmetry breaking pattern were described by Cvetic *et al.* (1992), and limits on the charged sector masses and mixings for general models without left-right symmetry have been given by Langacker and Uma Sankar (1989) and Yao *et al.* (2006).

The simple forms of the (supersymmetric) left-right model are not consistent with gauge unification unless the $SU(2)_R$ breaking occurs at a much higher scale (e.g., 10^{12} GeV). Such a large scale is also required by current allowed ranges for the neutrino masses in the triplet versions. In some cases, the initial breaking can leave $U(1)_{3R}\times U(1)_{BL}$ unbroken. Realistic SO(10) breaking patterns suggest α in the range 0.7–0.9 (Robinett and Rosner, 1982b). An important special case is the χ model, which occurs when SO(10) breaks directly to $SU(5)\times U(1)_\chi$. This corresponds to Eq. (42) for $\kappa=1$ and $\sin^2\theta_W=3/8$ [which is the value predicted by SU(5) at the *unification scale*], leading to $\alpha=\sqrt{2/3}\sim 0.82$.

A generalization of this type of model is based on $SU(2)\times U(1)_Y\times U(1)_2$, where Q_2 is a linear combination

$$Q^{YBL} = aY + bT_{BL} \equiv b(zY + T_{BL}), \tag{45}$$

where $b\neq 0$. It is convenient to normalize b so that the coupling g_2 is given by Eq. (44), or alternatively one can choose $b=1$ and take g_2 to be arbitrary. The $U(1)_{3R}\times U(1)_{BL}$ limit in Eq. (42) corresponds to choosing $b^2 z(1+z) = -3/5$ with $\alpha\equiv\sqrt{5/3}bz$. Q^{YBL} is anomaly free for the standard model fermions (including ν_R) (Weinberg, 1996). Y and Q^{YBL} are nonorthogonal (i.e., $\Sigma_f Y_f Q_f^{YBL}\neq 0$ when summed over a family of the known left-handed fermions and antifermions), except for the special case of $U(1)_{3R}\times U(1)_{BL}$, but it could come about by a more general embedding of the generators or by (possibly large) kinetic mixing, as discussed in Sec. II.C. The pure B-L model ($z=0$) is often studied phenomenologically, and has the property that the ordinary Higgs doublets do not induce Z-Z' mixing. The models in this class have been systematically discussed by Appelquist *et al.* (2003), including generalizations with an arbitrary number of ν_R with nonuniversal charges.

This entire class of models based on T_{3R} and T_{BL} (or Y and T_{BL}) is perhaps less interesting in a supersymmetric context because the two supersymmetric Higgs doublets $H_{u,d}$ form a vector pair with $T_{3R}=\pm\frac{1}{2}$ and $T_{BL}=0$. Therefore, an elementary μ term in Eq. (39) is not forbidden by the extra $U(1)'$. Similar difficulties apply to the SM singlet supermultiplets that are needed to break the $U(1)'$ since they would most likely be introduced as nonchiral vector pairs to avoid anomalies. (One could instead give large VEVs to the scalar partners of the ν^c, but this would break R parity and would be challenging for neutrino phenomenology.)

3. The E_6 models

Many Z' studies focus on the two extra $U(1)'$'s which occur in the decomposition of the E_6 GUT (Robinett and Rosner, 1982a; Langacker *et al.*, 1984; Hewett and Rizzo, 1989), i.e., $E_6\to SO(10)\times U(1)_\psi$ and $SO(10)\to SU(5)\times U(1)_\chi$. We consider them only as simple examples of anomaly-free $U(1)'$ charges and exotic fields, and do not assume a full underlying grand unified theory. In E_6, each family of left-handed fermions is promoted to a fundamental **27**-plet, which decomposes under $E_6\to SO(10)\to SU(5)$ as

$$\mathbf{27}\to\mathbf{16}+\mathbf{10}+\mathbf{1}\to(\mathbf{10}+\mathbf{5}^*+\mathbf{1})+(\mathbf{5}+\mathbf{5}^*)+\mathbf{1}, \tag{46}$$

as shown in Table II. In addition to the standard model fermions, each **27**-plet contains two standard model singlets ν^c and S [which may be charged under the $U(1)'$]. ν^c may be interpreted as the conjugate of the right-handed neutrino. There is also an exotic color-triplet quark D with charge $-1/3$ and its conjugate D^c, both of which are SU(2) singlets, and a pair of color-singlet SU(2)-doublet exotics, $H_u=\binom{H_u^+}{H_u^0}$ and $H_d=\binom{H_d^0}{H_d^-}$ with $y_{H_{u,d}}=\pm 1/2$. H_d transforms the same way as $H_u^c\equiv\tilde{H}_u$, the (tilde) conjugate of H_u under the SM. The exotic fields are all therefore singlets or nonchiral under the standard model, but may be chiral under the $U(1)'$.

The E_6 models can be considered in both nonsupersymmetric and supersymmetric versions. In the supersymmetric case, the scalar partners of the S and ν^c can develop VEVs to break the $U(1)'$ symmetry, though the latter (as well as a VEV for the scalar partner of the ν) would break R parity and may be problematic for neutrino phenomenology. Similarly, the scalar partners of one $H_{u,d}$ pair can be interpreted as the two Higgs doublets of the MSSM. The two additional $H_{u,d}$ families may be interpreted either as additional Higgs pairs or as exotic leptons (H_d has the same SM quantum numbers as an ordinary lepton doublet, while H_u would be conjugate to a right-handed exotic doublet).

Table II also lists the $U(1)_\chi$ and $U(1)_\psi$ charges of the **27**-plet. By construction, the fields in an irreducible representation of SO(10) [SU(5)] all carry the same ψ (χ) charges. Most studies assume that only one Z', coupling to the linear combination

$$Q(\theta_{E_6}) = \cos\theta_{E_6} Q_\chi + \sin\theta_{E_6} Q_\psi, \qquad (47)$$

where $0 \leqslant \theta_{E_6} < \pi$ is a mixing angle, is relevant at low energies. [One can also include a kinetic mixing correction $-\epsilon Y$ to the effective charge, as in Eq. (34).] As discussed in Sec. III.A.2, the χ model ($\theta_{E_6}=0$) is a special case of the T_{3R} and B-L models, supplemented with additional exotic fields in the **10**+**1** of SO(10). Since the latter are nonchiral in this case they may be omitted, or one or more **10**'s may be introduced as Higgs fields. The ψ model ($\theta_{E_6}=\pi/2$), on the other hand, has chiral exotics and requires the three full **27**-plets. Using Eq. (11) one sees that the currents of the fields in the **16** and **10** have purely axial couplings to the Z_ψ (this only holds for ν if it pairs with ν^c to form a Dirac fermion). Another commonly studied case is the η model,

$$Q_\eta = \sqrt{\frac{3}{8}}Q_\chi - \sqrt{\frac{5}{8}}Q_\psi$$
$$= -Q(\theta_{E_6} = \pi - \arctan\sqrt{5/3} \sim 0.71\pi),$$

which occurs in Calabi-Yau compactifications of the heterotic string if E_6 breaks directly to a rank 5 group (Witten, 1985) via the Wilson line (Hosotani) mechanism. The inert model $Q_I = -Q(\theta_{E_6}=\arctan\sqrt{3/5}\sim 0.21\pi)$ has a charge orthogonal to Q_η and follows from an alternative E_6 breaking pattern (Robinett and Rosner, 1982a). In the neutral N model ($\theta_{E_6}=\arctan\sqrt{15}\sim 0.42\pi$) (Ma, 1996; Barger *et al.*, 2003; Kang, Langacker, and Li, 2005; King *et al.*, 2006), ν^c has zero charge, allowing a large Majorana mass or avoiding big bang nucleosynthesis constraints for a Dirac ν, as discussed in Sec. V.F. It essentially interchanges the assignments of S and ν^c and of the two $\mathbf{5}^*$ representations (which have the same standard model quantum numbers) with respect to the χ model, and is basically the same as the alternative left-right model by Babu *et al.* (1987) and Ma (1987). The secluded sector model [$\theta_{E_6}=\arctan(\sqrt{15}/9)\sim 0.13\pi$] (Erler *et al.*, 2002) will be discussed in Sec. III.E.3.

The E_6 models allow the Yukawa couplings needed to generate masses for the standard model and exotic fermions. In particular, in the supersymmetric case the superpotential terms

$$W = -h_d Q H_d d^c + h_u Q H_u u^c - h_e L H_d e^+ + h_\nu L H_u \nu^c$$
$$+ \lambda_S S H_u H_d + \lambda_D S D D^c \qquad (48)$$

are all allowed, where family indices have been neglected. (In the non-SUSY case, two Higgs doublets, analogous to H_u and H_d, are required.) From Eq. (48) we see that the E_6 models all allow a dynamical μ_{eff}, while an elementary μ is forbidden in all but the χ model.

The supersymmetric E_6 model with three **27**-plets can incorporate one or more pairs of Higgs doublets $H_{u,d}$ in the $\mathbf{5}+\mathbf{5}^*$ pairs. However, that version of the model is not consistent with the simple form of gauge unification observed in the MSSM for the SM subgroup. That is because the complete extra $\mathbf{5}+\mathbf{5}^*$ multiplets give equal contributions to the SU(3), SU(2), and $U(1)_Y$ β functions at one loop, so the unification conditions are similar to the MSSM with three families but no Higgs pair. Unification can be restored by introducing an H_u and H_u^c pair from an incomplete $\mathbf{27}+\mathbf{27}^*$ representation (Langacker and Wang, 1998). (The physical H_u could either be this one or from the complete **27**-plets.) This pair is completely nonchiral so it does not introduce any anomalies, but at the cost of introducing a rather arbitrary aspect to the model. Also, there is no obvious reason [except perhaps the mechanism in Giudice and Masiero (1988)] why this extra pair should be at the electroweak or TeV scale, reintroducing a form of the μ problem. Nevertheless, the unification of the SM gauge couplings and the unification scale M_X are then the same as in the MSSM at one loop, though the value of the gauge coupling at M_X is increased because of the extra exotics.

If the $U(1)'$ really derives from an E_6-type GUT which breaks directly to $SU(3)\times SU(2)\times U(1)_Y\times U(1)'$, one expects that $g_2=\sqrt{\frac{5}{3}}g'$ at the unification scale, where $\sqrt{\frac{5}{3}}g'$ is the GUT-normalized hypercharge coupling. Running down to the TeV scale, this implies

$$g_2 = \sqrt{\frac{5}{3}} g \tan\theta_W \lambda_g^{1/2}, \qquad (49)$$

where $\lambda_g^{1/2}\sim 1$ up to a (θ_{E_6}-dependent) correction of a few % due to the $U(1)'$ charge of the incomplete **27** $+\mathbf{27}^*$. Equation (49) can be taken as a definition of λ_g for an arbitrary model. It is typically of order unity even for more complicated E_6 breaking patterns (Robinett and Rosner, 1982a), and was taken to be unity by construction for the G_{LR} model.

In a full E_6 grand unified theory the exotic D, D^c partners of the Higgs doublets would have diquark Yukawa couplings such as $W_{DQ}\sim DQQ$ or $D^c u^c d^c$, as well as leptoquark couplings $W_{LQ}\sim D u^c e^c$ or $D^c QL$, which are related by E_6 to the ordinary Higgs Yukawa couplings. These would lead to rapid proton decay mediated by the D and D^c unless their masses [and therefore the $U(1)'$ breaking scale] is comparable to the unification scale. A TeV-scale Z' therefore requires that the GUT Yukawa relations are not respected so that either the leptoquark or diquark couplings (or both) are absent. This could come about in a string construction if the fields in the multiplet are not directly related to each other in the underlying theory [see, e.g., Witten (2001)]. See King *et al.* (2006) and Howl and King (2008) for a detailed study of complete E_6 models with a low energy $U(1)'$. Alternatively, one can simply view the charges and exotics as an example of an anomaly-free construction.

B. Anomaly-free sets

Many authors have described classes of $U(1)'$ models by requiring the cancellation of anomalies and other criteria (Barr *et al.*, 1986; Cvetic *et al.*, 1997; Cheng *et al.*, 1998, 1999; Erler, 2000; Joshipura *et al.*, 2000; Ma, 2002; Appelquist *et al.*, 2003; Carena *et al.*, 2004; Demir *et al.*, 2005; Batra *et al.*, 2006; Morrissey and Wells, 2006; Kang *et al.*, 2008; Langacker *et al.*, 2008; Lee *et al.*, 2008). Usually some conditions are applied on the types of exotics. It is usually assumed that any exotic fermions are nonchiral under the standard model, i.e., that they occur in vector pairs $\psi+\psi^c$. This avoids the introduction of any SM anomalies and also reduces the sensitivity to precision electroweak constraints (Yao *et al.*, 2006). One can then constrain the exotic representations with respect to the SM and their $U(1)'$ charges from the mixed SM-$U(1)'$ conditions

$$\sum_{f\in 3,3^*} Q_{2f}=\sum_{f\in 2} Q_{2f}=\sum_{f} Y^2 Q_{2f}=\sum_{f} Y Q_{2f}^2=0. \qquad (50)$$

The pure $U(1)'$ conditions $\Sigma_f Q_{2f}=\Sigma_f Q_{2f}^3=0$ further restrict the charges. Alternatively, some ignore the latter because they can be satisfied by adding SM singlets to the model. This can always be done with rational charges[4] if the mixed anomaly solutions are rational (Batra *et al.*, 2006; Morrissey and Wells, 2006). However, because of its cubic nature the singlet structure is sometimes complicated.

In nonsupersymmetric models it is often assumed that the only chiral fermions are the three ordinary families and three corresponding families of exotics. This assumption is often not valid in supersymmetric models, where one must also take into account the fermionic partners of the Higgs doublets and SM singlets which break the $U(1)'$. These, as well as other exotics, often do not occur in three families (exceptions are the E_6 models, where they do occur in three families, and the T_{3R}, B-L models, where the Higgs doublets and singlets are usually nonchiral).

Another issue in the supersymmetric models is whether the MSSM unification of the SM gauge couplings is preserved. The simplest way for this to occur is for the three SM families, which transform as $\mathbf{10}+\mathbf{5}^*$ under SU(5), and the two Higgs doublets $H_{u,d}$, are supplemented by exotics which transform as $\mathbf{5}+\mathbf{5}^*$ and/or $\mathbf{10}+\mathbf{10}^*$. It is not necessary for the fields in a SU(5) multiplet to have the same $U(1)'$ charges (e.g., they may have different origins in an underlying string theory), and in fact under minimal assumptions they must be different (Morrissey and Wells, 2006). An alternative is to allow nonchiral exotics, as in the E_6 models, reintroducing a form of the μ problem.

Other conditions are often employed along with the anomaly and unification constraints. These may involve the existence of quark and lepton Yukawa couplings for one or two Higgs doublets, constraints on neutrino mass, Yukawa couplings that can lead to masses for the exotics, operators that can allow exotic decays, whether the charges are family universal, whether the $U(1)'$ solves the supersymmetric μ problem, whether it forbids R-parity violating operators (Erler, 2000; Joshipura *et al.*, 2000; Ma, 2002) or other operators relevant to proton decay (Chamseddine and Dreiner, 1995; Lee, Luhn, and Matchev, 2007; Coriano *et al.*, 2008), whether it plays the role of a family symmetry relevant to the fermion masses and mixings (Kaplan *et al.*, 1999; Joshipura *et al.*, 2000), and many other possible conditions.

The Q^{YBL} models of Sec. III.A.2, which do not require any exotics other than ν_R, are discussed in Appelquist *et al.* (2003). Four one-parameter families of models with three families of exotics were constructed in Carena *et al.* (2004). Two of these, referred to as $q+xu$ and $\mathbf{10}+x\mathbf{5}^*$, are equivalent to the Q^{YBL} and the E_6 model $Q(\theta_{E_6})$, respectively, while the others (B-xL and d-xu) have not emerged from other considerations for general x. The $\mathbf{10}+x\mathbf{5}^*$ and d-xu would require two Higgs doublets to have normal quark and charged lepton Yukawas.

The most systematic classification of the supersymmetric models is given in Erler (2000), which requires anomaly cancellation, minimal gauge unification with no nonchiral states, exotic masses, and the absence of rapid proton decay or fractional electric charges. Classes of solutions were found, which required that more than one SM singlet participates in the $U(1)'$ breaking. A particularly simple one is the $Q_{\tilde\psi}$ model. It involves two $\mathbf{5}+\mathbf{5}^*$ pairs D_i+L_i and $D_i^c+L_i^c$, $i=1,2$, which are analogous to the (D,H_u) and (D^c,H_d) of the E_6 model, along with $H_{u,d}$ and the three SM families. The $U(1)'$ symmetry is broken by the VEVs of two singlets S and S_D, which also generate masses for the $H_{u,d}$ and $L_i, L_i^c(\langle S\rangle)$ and for $D_i, D_i^c(\langle S_D\rangle)$. Additional singlets are needed for the $U(1)'$ anomalies. The $Q_{\tilde\psi}$ charges are listed in Table III. The fermion currents are purely axial. It is straightforward to generalize the $Q_{\tilde\psi}$ model to $Q_{\mathbf{55}^*}$, which allows nonaxial charges and $n_{\mathbf{55}^*}$ pairs of $\mathbf{5}+\mathbf{5}^*$. Three distinct chiral singlets must acquire VEVs to generate all of the exotic masses, except for $n_{\mathbf{55}^*}=2$ or 3. Additional SM singlets are needed for the $U(1)'$ anomalies and to generate singlet masses. The gauge coupling g_2 is arbitrary.

It was shown in Demir *et al.* (2005) that anomaly-free supersymmetric models can be constructed without any exotics (not even ν^c) and only one singlet S (which generates a dynamical $\mu_{\rm eff}$) provided one allows family non-

[4]One expects the charges to be rational if the $U(1)'$ is embedded in a simple group, but this need not be the case for more complicated embeddings, such as the SM couplings in Eq. (8).

TABLE III. Examples of supersymmetric models consistent with minimal SM gauge unification. n_{55^*} is the number of pairs of $\mathbf{5}+\mathbf{5}^*$. Q_S is taken to be 1. The free parameters are $Q_{H_u}\equiv x, Q_Q\equiv y, Q_D\equiv z$ (which only affects the exotics), and the gauge coupling g_2. Kinetic mixing can be added. The $Q_{\tilde{\psi}}$ model is a special case with axial charges and $n_{55^*}=2$. Additional SM singlets are not displayed. The ν^c charge allows a Dirac ν mass term.

	Q_{55^*}	$Q_{\tilde{\psi}}$		Q_{55^*}	$Q_{\tilde{\psi}}$
Q	y	1/4	H_u	x	$-1/2$
u^c	$-x-y$	1/4	H_d	$-1-x$	$-1/2$
d^c	$1+x-y$	1/4	S_D	$3/n_{55^*}$	3/2
L	$1-3y$	1/4	D_i	z	$-3/4$
e^+	$x+3y$	1/4	D_i^c	$-3/n_{55^*}-z$	$-3/4$
ν^c	$-1-x+3y$	1/4	S_L	$2/n_{55^*}$	1
S	1	1	L_i	$\frac{5-n_{55^*}}{4n_{55^*}}+x+3y+3z/2$	$-1/2$
			L_i^c	$-2/n_{55^*}-Q_{L_i}$	$-1/2$

universal charges [an early example was also given in Cvetic *et al.* (1997)]. It is possible to choose the charges to avoid flavor changing neutral current (FCNC) effects (see Sec. V.D). However, the U(1)′ forbids some of the quark and lepton Yukawa interactions in the superpotential. These could possibly be generated by nonholomorphic soft terms, as described in Sec. II.E.

C. TeV scale physics models

In this section we consider various models involving new TeV scale physics, especially those motivated as alternatives to the elementary Higgs for electroweak symmetry breaking.

As a preliminary, consider a direct product of two identical gauge group $G\equiv G_1\times G_2$, with generators $\vec{T}_{1,2}$ and associated currents $\vec{J}_{1,2}$. Then

$$-L=g_1\vec{J}_1\cdot\vec{W}_1+g_2\vec{J}_2\cdot\vec{W}_2. \tag{51}$$

G can be spontaneously broken to the *diagonal subgroup* G_D with generators $\vec{T}_D=\vec{T}_1+\vec{T}_2$ if there is a Higgs field which transforms equivalently under both groups. An example is $SU(N)\times SU(N)$, with a Higgs φ_a^α transforming as $\mathbf{N}^*\times\mathbf{N}$, with $\langle\varphi_a^\alpha\rangle=c\delta_a^\alpha$. It is then straightforward to show that

$$-L=g_L(\vec{J}_1+\vec{J}_2)\cdot\vec{W}_L+g_L(\cot\delta\vec{J}_1-\tan\delta\vec{J}_2)\cdot\vec{W}_H, \tag{52}$$

where $\vec{W}_L=\sin\delta\vec{W}_1+\cos\delta\vec{W}_2$ is the massless boson, W_H is the massive orthogonal combination, $\tan\delta=g_2/g_1$, and $g_L=g_1\sin\delta$. W_L can acquire mass and $W_{L,H}$ can mix due to additional Higgs fields. A simple illustration is the SM breaking of $U(1)_{T_3}\times U(1)_Y$ to $U(1)_{em}$ by the ordinary Higgs doublet.

1. Little Higgs, twin Higgs, and ununified models

In *little Higgs* models (Arkani-Hamed, Cohen, and Georgi, 2001) the Higgs is a pseudo-Goldstone boson of an approximate global symmetry. [For reviews, see Han *et al.* (2003, 2006) and Perelstein (2007).] The one-loop (and sometimes two-loop) quadratic divergences in the Higgs mass square are canceled by new TeV gauge bosons, fermions, and scalar particles related to those of the SM. There are a wide class of models, all of which involve heavy neutral and charged gauge bosons. For example, in the littlest Higgs models (Arkani-Hamed *et al.*, 2002) the electroweak gauge group is $[SU(2)\times U(1)]^2$, which is a subgroup of a larger global group. The SM left-handed fermions are charged under only the first SU(2). The $SU(2)^2$ symmetry is broken by a condensate charged under both factors to an unbroken diagonal subgroup, and the U(1) charges are chosen to yield $U(1)_Y\times U(1)_H$, where Y is the normal hypercharge. Thus, the residual gauge group is $SU(2)_L\times U(1)_Y\times SU(2)_H\times U(1)_H$. From Eq. (52), the heavy charged $W_H^\pm$ and neutral W_H^0 couple to the left-handed SM quarks and leptons with the $SU(2)_L$ generators $\vec{\tau}/2$ and with coupling $g\cot\delta$. The neutral $U(1)_H$ boson is lighter, with model dependent couplings. Precision electroweak constraints are rather severe, unless one pushes the little Higgs scale to be uncomfortably large compared to the original motivation. However, the difficulties can be reduced if the $U(1)_H$ is not gauged.

The precision electroweak constraints are greatly weakened (they are only generated at loop level) if one introduces a discrete symmetry, T parity (Cheng and Low, 2003; Hubisz *et al.*, 2006). This is analogous to R parity in supersymmetry, and requires that the heavy states, such as the new gauge bosons, only couple in pairs to the ordinary particles. This also means that they must be pair produced at colliders. The lightest could be stable and possibly be a dark matter candidate. However, it has recently been argued that the T parity may be broken by anomalies (Hill and Hill, 2007), leading to decays, for example, into ZZ.

In the *twin Higgs* model (Chacko *et al.*, 2006) the Higgs quadratic divergences are canceled by particles from a hidden sector that is a mirror of the SM and which mainly communicates by an extended Higgs sector. The gauge bosons in the hidden sector may essentially decouple from the SM particles and could even be massless, while in other versions there may be kinetic mixing with the photon.

In the *ununified* model (Georgi *et al.*, 1990), the left-chiral SM quarks and leptons transform under distinct SU(2) groups $SU(2)_q$ and $SU(2)_l$ with gauge couplings $g_{q,l}$, i.e., they are not unified. There is a single conventional $U(1)_Y$. After diagonal breaking, one recovers the SM along with heavy $W_H^{\pm,0}$ which couple to $g(\cot \delta \vec{J}_q - \tan \delta \vec{J}_l)$ using Eq. (52). For small $\tan \delta$ the heavy bosons couple mainly to quarks.

2. Extra dimensions

The existence of extra dimensions is suggested by string models (Antoniadis, 1990). There are a wide variety of models, depending on their number, size, whether they are flat or warped, whether the SM fields are allowed to propagate in the extra dimensions (i.e., in the *bulk*), etc. For a review, see Yao *et al.* (2006).

The simplest case involves a single extra dimension of radius R, implying the existence of Kaluza-Klein excitations of the states that can propagate in the bulk, with mass $\sim n/R, n=1,\ldots$. If only gravitons propagate, then R can be large enough to probe in laboratory gravity experiments. However, if the SM gauge bosons are also allowed to propagate, then R^{-1} must be larger than $O(\text{TeV})$ ($R \lesssim 10^{-17}$ cm). If the SM fermions and Higgs are not allowed to propagate (i.e., confined to the brane), then the excitations of the SM gauge bosons ($W^{\pm}, Z, A$, gluon) couple to the same currents as their SM counterparts, but with a coupling constant larger by $\sqrt{2}$ (Casalbuoni *et al.*, 1999; Masip and Pomarol, 1999). Current experimental limits require $R^{-1} \gtrsim 7$ TeV (Cheung and Landsberg, 2002; Barbieri *et al.*, 2004). The limits are much weaker [$O(300\text{ GeV})$] in universal extra dimension models, in which all of the SM fields propagate uniformly in the bulk (Appelquist *et al.*, 2001; Cheng *et al.*, 2002; Appelquist and Yee, 2003; Gogoladze and Macessnu, 2006). Similar to R or T parity, there is a Kalvza-Klein parity so that the $n=1$ states can only be pair produced and only contribute to electroweak observables in loops. The lightest is stable. In variants in which the various quarks and leptons are localized in different parts of the extra-dimensional space (with implications for the flavor problem) the couplings of the Kaluza-Klein excitations are family nonuniversal (since the overlap of the wave functions depends on location). This leads to the possibility of FCNC effects (Delgado *et al.*, 2000), as discussed in Sec. V.D.

Models involving warped extra dimensions (Randall and Sundrum, 1999) may have all of the SM fields confined to the infrared brane. However, much attention has been devoted to the possibility that the SM fields other than the Higgs can also propagate in the bulk (Hewett *et al.*, 2002; Agashe *et al.*, 2003, 2007; Carena, Delgado, *et al.*, 2003a), because in that case the theory is related to technicolor models by the anti–de Sitter–conformal field theory correspondence (Maldacena, 1998). It is then useful to enhance the electroweak gauge symmetry to $SU(2)_L \times SU(2)_R \times U(1)_{B\text{-}L}$ to provide a custodial symmetry to protect the electroweak ρ parameter (Agashe *et al.*, 2003). The Kaluza-Klein excitations of the gauge bosons couple mainly to the t and b due to wave function overlaps, and decays to WW and $Z+$ Higgs are also possible due to mixings [see, e.g., Agashe *et al.* (2007)].

3. Strong dynamics

There have been many models in which strong dynamics is involved in electroweak symmetry breaking, which often involve additional elementary gauge bosons or composite spin-1 states, which may be strongly coupled.

Dynamical symmetry breaking (DSB) models in which the Higgs is replaced by a fermion condensate are reviewed in Chivukula and Simmons (2002), Hill and Simmons (2003), and Chivukula *et al.* (2004). For example, *top-color* models (Hill, 1995) typically involve new gluons and a new Z' that couple preferentially and with enhanced strength to the third generation and which assist in forming a top condensate. Nonuniversal extended technicolor models (Chivukula and Simmons, 2002) also feature new gauge interactions preferentially coupled to the third family.

The breaking electroweak symmetry strongly (BESS) models (Casalbuonl *et al.*, 1985, 1987) are effective Lagrangian descriptions of models with a strongly interacting longitudinal gauge boson sector, such as one expects in the large M_H limit of the SM or in some forms of DSB. There are vector and axial bound states which can mix with the $W^{\pm}$, Z, and A. They interact with the SM particles directly and by mixing. The possibility that the electroweak bosons could be composite has also been considered (Baur *et al.*, 1987).

Another interesting model with no elementary or composite Higgs (Csaki *et al.*, 2004) is a variant on the warped extra dimension scenario. However, instead of including a Higgs field the electroweak symmetry is broken by boundary conditions. The Kaluza-Klein excitations of the gauge bosons unitarize the high-energy scattering of longitudinal gauge bosons. More general classes of Higgless models may involve fermiophobic Z' which may be produced and detected by their couplings to the W and Z (He *et al.*, 2008).

D. Nonstandard couplings

Most of the canonical Z' models assume electroweak scale couplings, and that the Z' couplings to most or all of the SM fermions are of comparable strength and family universal, in which case existing experimental constraints require masses not too much below 1 TeV (Sec. IV). However, there are many models with different assumptions concerning the gauge couplings, charges, and scales.

1. Decoupled models

Leptophobic Z''s (del Aguila, Blair, *et al.*, 1987) do not couple to ordinary neutrinos or charged leptons, and therefore most direct electroweak and collider searches are insensitive to them. They could emerge in the E_6 η model in Table II, combined with a (large) kinetic mixing (Babu *et al.*, 1996) $\epsilon \sim -1/\sqrt{15}$; in a flipped SU(5) model (Lopez and Nanopoulos, 1997); or in models in which the Z' couples to baryon number (Carone and Murayama, 1995). Approximately leptophobic models were once suggested by apparent anomalies in $Z \to b\bar{b}$ decays [see Rosner (1996b) and Umeda *et al.* (1998) for references], but are still an interesting possibility for allowing Z' masses much smaller than a TeV. A purely leptophobic Z' is still constrained by Z-Z' mixing effects (Umeda *et al.*, 1998), and could be inferred by collider signals such as the production of $t\bar{t}$ pairs, exotics (Rosner, 1996b), or the same-sign dilepton decays of a pair of heavy Majorana neutrinos (Duncan and Langacker, 1986; del Aguila and Aguilar-Saavedra, 2007). They could even be light enough to be produced in Υ decays (Aranda and Carone, 1998).

Limits are also weak if the Z' couples only to the second and third family leptons (Foot *et al.*, 1994) or if it couples only to third family fermions (Andrianov *et al.*, 1998). In *fermiophobic* models (Barger *et al.*, 1980; Donini *et al.*, 1997) there are no direct couplings of the Z' to the SM fermions, although they may be induced by ordinary or kinetic mixing. An interesting possibility is that such fermiophobic Z' may couple to a hidden sector (Chang *et al.*, 2006; Kumar and Wells, 2006), such as may be associated with supersymmetry breaking. Mixing effects could therefore possibly be a means of probing such a sector (direct Z' couplings to a hidden sector are considered in Sec. V.E). Finally, a Z' with canonical charges could still be much lighter than a TeV if its gauge coupling is sufficiently small (Fayet, 1980, 2007; Freitas, 2004; Nelson and Walsh, 2008).

2. Stückelberg models

It is possible to write a U(1) gauge invariant theory with a massive gauge boson C_μ by the Stückelberg mechanism (Stueckelberg, 1938). The Lagrangian is

$$L = -\tfrac{1}{4}C^{\mu\nu}C_{\mu\nu} - \tfrac{1}{2}(mC^\mu + \partial^\mu\sigma)(mC_\mu + \partial_\mu\sigma), \tag{53}$$

where $C_{\mu\nu}$ is the field strength tensor. Under a gauge transformation, $\Delta C_\mu = \partial_\mu\beta$, while the field σ is shifted, $\Delta\sigma = -m\beta$, analogous to the shift in an axion field. A gauge-fixing term can be added to Eq. (53) which cancels the cross term between C and σ, leaving a massive C field and a decoupled σ. This is analogous to the Higgs mechanism, but there is no field with a VEV and no physical Higgs boson. This mechanism has recently been applied to a U(1)$'$ extension of the SM or the MSSM (Kors and Nath, 2004; Feldman *et al.*, 2007). For example, if one replaces the second term in Eq. (53) by $-(M_2C_\mu + M_1B_\mu + \partial_\mu\sigma)^2/2$ with $M_1/M_2 \equiv \epsilon \ll 1$, then C will mix with A and Z, but there will remain a massless photon. The new Z_2 can be relatively light (e.g., several hundred GeV), so ϵ must be small. If the C has no direct couplings to matter, the new Z_2 will decay only to SM particles via the mixing and will be very narrow. If C does couple to exotic matter, then the mixing with the photon will induce tiny (generally) irrational electric charges of $O(\epsilon)$ for the exotic particles. Such mixing with the photon is never induced by ordinary Higgs-type mixing if $U(1)_{em}$ is unbroken, but can also be induced by kinetic mixing with another massless boson (Sec. III.E.1). Other applications, such as to dark matter, are reviewed in Feldman *et al.* (2007).

3. Family nonuniversal models

Another variant is the possibility of *family nonuniversal* charges [see, e.g., Demir *et al.* (2005)]. A number of examples of Z' coupling preferentially to the third family or to the t quark were mentioned in Secs. III.C.2, III.C.3, and III.D.1. These could have enhanced gauge couplings, and could be observed as a resonance in $t\bar{t}$ production. String-derived Z''s often have nonuniversal couplings as well (Sec. III.F), as do the Kaluza-Klein excitations in extra-dimensional theories in which the fermion families are spatially separated (Sec. III.C.2). Possible FCNC effects are considered in Sec. V.D.

E. U(1)$'$ breaking scales

Most attention is given to possible electroweak or TeV scale Z''s, but there are other possibilities. Here we describe massless, TeV scale, and intermediate scale models. Models involving the GUT or string scales are described in Sec. III.F.

1. Massless Z'

A Z' would be massless if the U(1)$'$ symmetry is unbroken. This would imply an unacceptable long range force if it coupled to ordinary matter unless the coupling were incredibly small (Dobrescu and Mocioiu, 2006). It would be allowed if the primary coupling were to a hidden sector and communicated only by higher-dimensional operators (Dobrescu, 2005) or by kinetic mixing with the photon (Holdom, 1986). The latter scenario would induce a small fractional electric charge for hidden sector particles.

2. Electroweak or TeV scale Z'

Models in which the U(1)$'$ is involved in electroweak symmetry breaking, such as in Sec. III.C, typically involve U(1)$'$ breaking at the electroweak or TeV scale.

In the U(1)′ extension of the MSSM with a single S field (Cvetic and Langacker, 1996a, 1997; Cvetic *et al.*, 1997; Keith and Ma, 1997; Langacker and Wang, 1998), the part of the superpotential involving S and $H_{u,d}$ is $W=\lambda_S S H_u H_d$, where we have assumed $Q_S \neq 0$ and $Q_S + Q_u + Q_d = 0$. Like the MSSM, the minimum of the tree-level potential always occurs along the charge-conserving direction with only $\langle H^0_{u,d}\rangle \neq 0$ (this assumes that the squark and slepton VEVs vanish). The potential is then

$$V = V_F + V_D + V_{\text{soft}}, \tag{54}$$

where

$$V_F = \lambda_S^2(|H_u^0|^2|H_d^0|^2 + |S|^2|H_u^0|^2 + |S|^2|H_d^0|^2),$$

$$V_D = \frac{g_1^2}{8}(|H_u^0|^2 - |H_d^0|^2)^2 + \frac{g_2^2}{2}(Q_u|H_u^0|^2 + Q_d|H_d^0|^2 + Q_S|S|^2)^2,$$

$$V_{\text{soft}} = m_u^2|H_u^0|^2 + m_d^2|H_d^0|^2 + m_S^2|S|^2 - (\lambda_S A_S S H_u^0 H_d^0 + \text{H.c.}). \tag{55}$$

If S acquires a VEV, then the effective μ parameter is $\mu_{\text{eff}}=\lambda_S\langle S\rangle$, the corresponding effective $B\mu$ is $(B\mu)_{\text{eff}} = \lambda_S A_S \langle S\rangle$, and the Z-Z' mass matrix is given by Eqs. (18) and (20). One can define the fields so that $\lambda_S A_S$ and therefore the VEVs $\nu_{u,d}$ and s defined after Eq. (20) are real and positive. There is no analog of the first (second) term in V_F (V_D) in the MSSM.

For generic parameters one expects $\nu_{u,d}$ and s to be comparable. For example, for $\lambda_S A_S$ large compared to the soft masses and $Q_u = Q_d = -Q_S/2$ one finds (Cvetic *et al.*, 1997) $\nu_u \sim \nu_d \sim s$, with negligible Z-Z' mixing and $M^2_{Z'}/M^2_Z \sim 12 g_2^2 Q_u^2/g_1^2$, which is typically of order 1. This case is excluded unless the model is leptophobic or something similar.

A more likely scenario is that the soft parameters ($|m_{u,d,S}|, |A_S|$) are of O(1 Tev), with $m_S^2 < 0$. Then $s^2 \sim -2m_S^2/g_2^2 Q_S^2$ and $M^2_{Z'} \sim -2m_S^2$. One can have a smaller electroweak scale $\nu_{u,d} \ll s$ by accidental cancellations, which are not excessive provided $M_{Z'}$ is not too much larger than a TeV. In most supersymmetry mediation schemes m_S^2 is positive at a large scale such as the Planck scale. The running m_S^2 can be driven negative at low scales radiatively provided it has sufficiently large Yukawa couplings, such as λ_S and/or couplings to exotics such as in Eq. (48). This is analogous to the MSSM in which m_u can be driven negative by its large Yukawa coupling to the top.

3. Secluded sector and intermediate scales

In the single S model in Eqs. (54) and (55) there is some tension between the electroweak scale and developing a large enough $M_{Z'}$. These can be decoupled without tuning when there are several S fields. For example, in the *secluded sector* model (Erler *et al.*, 2002) there are four standard model singlets S, $S_{1,2,3}$ that are charged under a U(1)′, with

$$W = \lambda_S S H_u H_d + \lambda S_1 S_2 S_3. \tag{56}$$

(Structures similar to this are often encountered in heterotic string constructions.) μ_{eff} is given by $\lambda_S\langle S\rangle$, but all four VEVs contribute to $M_{Z'}$. The only couplings between the ordinary $(S, H_{u,d})$ and secluded $(S_{1,2,3})$ sectors are from the U(1)′ D term and the soft masses [special values of the U(1)′ charges, which allow soft mixing terms, are required to avoid unwanted additional global symmetries]. It is straightforward to choose the soft parameters so that there is a runaway direction in the limit $\lambda \to 0$, for which the ordinary sector VEVs remain finite while the S_i VEVs become large. For λ finite but small, the $\langle S_i\rangle$ and $M_{Z'}$ scale as $1/\lambda$. For example, one can find $M_{Z'}$ in the TeV range for $\lambda \sim 0.05-0.1$. The secluded model can be embedded in the E_6 context (Table II).

Intermediate scale models (Cleaver *et al.*, 1998; Morrissey and Wells, 2007) are those in which the U(1)′ breaking is associated with a F and D flat direction, such as the secluded model in Eq. (56) with $\lambda = 0$. However, we consider a simpler toy model with two fields $S_{1,2}$ with $Q_{S_1} Q_{S_2} < 0$. If there are no terms in W like $S_i S_j$ or $S_i S_j S_k$, then the potential for $S_{1,2}$ is

$$V(S_1,S_2) = m_1^2|S_1|^2 + m_2^2|S_2|^2 + \frac{g_2^2}{2}(Q_{S_1}|S_1|^2 + Q_{S_2}|S_2|^2)^2. \tag{57}$$

The quartic term vanishes for $|S_2|^2/|S_1|^2 = -Q_{S_1}/Q_{S_2}$. For simplicity, take $Q_{S_1} = -Q_{S_2}$, and assume that at low energies $m^2_{S_1} < 0$ and $m^2_{S_2} > 0$, as would typically occur by the radiative mechanism if W contains a term $h_D S_1 D D^c$. If $m^2 \equiv m^2_{S_1} + m^2_{S_2} > 0$ the minimum will occur at $\langle S_1\rangle \neq 0, \langle S_2\rangle = 0$. If there is also a $\lambda_S S_1 H_u H_d$ term in W then $\langle S_1\rangle$ and $M_{Z'}$ will be at the electroweak scale ($\lesssim 1$ TeV), just as in the case of a single S. On the other hand, for $m^2 < 0$, the potential along the F and D flat direction $S_1 = S_2 \equiv S$ is

$$V(S) = m^2 S^2, \tag{58}$$

which appears to be unbounded from below. In fact, $V(S)$ is typically stabilized by one or both of two mechanisms: (a) The leading loop corrections to the effective (RGE-improved) potential result in $m^2 \to m^2(S)$, leading to a minimum slightly below the scale at which $m^2(S)$ goes through zero, which can be anywhere in the range 10^3-10^{17} GeV. (b) Another possibility is that the F flatness is lifted by higher-dimensional operators in W, such as $W = (S_1 S_2)^k/M^{2k-3}$, where M is the Planck or some other large scale. This would lead to $\langle S\rangle \sim \sqrt{mM} \sim 10^{11}$ GeV for $k=2$, $m \sim 1$ TeV, and M the Planck scale. In such models, higher-dimensional operators such as $L H_d e^+ (S/M)^p$ or $S H_u H_d (S/M)^q$ could also be important for generating small effective Yukawa couplings (and therefore fermion mass hierarchies) or $\mu_{\text{eff}} \ll \langle S\rangle$

terms (Cleaver *et al.*, 1998). Implications for neutrino mass are considered in Sec. V.F.

F. Grand unification, strings, and anomalous U(1)′

1. Grand unification

In some full grand unified theories (Langacker, 1981; Hewett and Rizzo, 1989), such as E_6, the extra U(1)′'s must be broken at or near the GUT unification scale to avoid rapid proton decay. This typically occurs if the Higgs doublets (and their Yukawa couplings to ordinary fermions) are related by the GUT symmetry to chiral exotics, which cannot be much heavier than the U(1)′ breaking scale. However, as mentioned in Sec. III.A.3 this can be evaded in models which respect the GUT quantum numbers but not the Yukawa relations, or in models such as the E_6 χ model, in which the Higgs doublets are nonchiral.

2. String theories

Most semirealistic superstring constructions yield effective four-dimensional field theories that include the SM gauge group (not a full four-dimensional GUT), as well as additional gauge group factors that often involve additional U(1)′'s. [Examples include Faraggi and Nanopoulos (1991), Faraggi (1993), Cleaver *et al.* (1999), Cvetic *et al.* (2001), Giedt (2001), Braun *et al.* (2005), Anastasopoulos, Dijkstra, *et al.* (2006), Coriano *et al.* (2008), and Lebedev *et al.* (2008). For reviews, see Cvetic and Langacker (1997) and Blumenhagen *et al.* (2005, 2007).] Heterotic constructions often descend through an underlying SO(10) or E_6 in the higher-dimensional space, and may therefore lead to the T_{3R} and B-L (i.e., Q^{LR}) or the E_6-type charges. Additional or alternative U(1)′ structures may emerge that do not have any GUT-type interpretation and therefore have very model-dependent charges. Similarly, intersecting brane constructions often descend through Pati-Salam type models (Mohapatra, 1986), yielding Q^{LR}. Other branes can lead to other types of U(1)′ charges. For example, the construction in Cvetic *et al.* (2001) involves two extra U(1)′'s, one coupling to Q^{LR} and the other only to the Higgs and the right-handed fermions.

Constructions often have one or multiple SM singlets which can acquire VEVs to break the extra U(1)′. However, that is not always the case. For example, in some of the Q^{LR} models [see, e.g., Cvetic *et al.* (2001) and Braun *et al.* (2005)] the only fields available to break the enhanced gauge symmetry are the scalar partners $\tilde{\nu}_R$ of the right-handed neutrinos (Cvetic *et al.*, 2002). These act like the δ_R^0 defined in Sec. III.A.2, but it is difficult to reconcile the Z' constraints with neutrino phenomenology. This also occurs in the simpler supersymmetric versions of the χ model.

The U(1)′ in string constructions may couple to hidden sector particles, and in some cases they can communicate between the ordinary and hidden sectors (Langacker *et al.*, 2008; Verlinde *et al.*, 2008). The nonstandard string U(1)′ often have family nonuniversal charges. This can occur if the fermion families have different embeddings in the underlying theory. A simple field-theoretic example is a variant on the E_6 model in Table II. One could assign the first two families (d_i^c, L_i), $i=1,2$, to the $\mathbf{5}^*$ from the $\mathbf{16}$ of SO(10), and the third to the $\mathbf{5}^*$ from the $\mathbf{10}$.

3. Anomalous U(1)′

The effective four-dimensional field theories arising from the compactification of a string theory usually contain anomalous U(1)′ factors [see Kiritsis (2004) for a review]. There is typically one anomalous combination in heterotic constructions. In intersecting brane models [Blumenhagen *et al.*, (2005)] there are stacks of branes yielding $U(N) \sim SU(N) \times U(1)$, in which the U(1) is usually anomalous. Since the underlying string theory is anomaly free, these anomalies must be canceled by a generalized Green-Schwarz mechanism. In particular, the Z' associated with the U(1)′ acquires a string-scale mass by what is essentially the Stückelberg mechanism in Eq. (53), with the axion field σ associated with an antisymmetric field in the internal space [this sometimes applies to nonanomalous U(1)′ as well]. The U(1)′ still acts as a global symmetry on the low-energy theory, restricting the possible couplings and having possible implications for baryon or lepton number. In addition, effective trilinear vertices may be generated between the Z' and the SM gauge bosons (Anastasopoulos, Bianchi, *et al.*, 2006a; Coriano *et al.*, 2006). It is possible that the string scale is actually very low (e.g., TeV scale) if there is a large total volume of the extra-dimensional space (a realization of the large extra dimension scenario). This would allow TeV scale Z''s associated with anomalous (or sometimes nonanomalous) U(1)′'s, without any associated Higgs scalar and with anomalous decays into ZZ, WW, and $Z\gamma$ (Ghilencea *et al.*, 2002; Berenstein and Pinansky, 2007; Armillis *et al.*, 2008; Kumar *et al.*, 2008).

Anomalous U(1)′'s in heterotic constructions lead to Fayet-Iliopoulos (FI) terms, which are effectively constant contributions to the U(1)′ D terms that are close to the string scale. Smaller FI terms may also appear in intersecting brane constructions which break supersymmetry. In many cases, FI terms trigger scalar fields in the low-energy theory to acquire VEVs to cancel them. These VEVs in turn may lead to the breaking of gauge symmetries [such as other nonanomalous U(1)′'s] and the generation of masses for some of the particles at the FI scale, a process known as vacuum restabilization [see Cleaver *et al.* (1999) for an example]. Family nonuniversal U(1)′'s may be used to generate fermion textures using the Froggatt-Nielsen mechanism (Ibanez and Ross, 1994; Binetruy and Ramond, 1995; Jain and Shrock, 1995; Chankowski *et al.*, 2005). The elements are associated with higher-dimensional operators allowed by the symmetry, and involve powers of the ratio of the FI and Planck scales.

IV. EXPERIMENTAL ISSUES

There are limits on Z' masses and Z-Z' mixing from precision electroweak data, from direct and indirect searches at the Fermilab Tevatron, and from interference effects at Large Electron Positron Collider 2 (LEP 2). In this section we review the existing limits and future prospects for discovery and diagnostics. FCNC effects for family nonuniversal couplings and astrophysical or cosmological constraints are discussed in Sec. V.

A. Constraints from precision electroweak

1. Parametrization

Precision electroweak data include purely weak νe and ν-hadron weak neutral current (WNC) scattering; weak-electromagnetic interferences in heavy atoms and in $e^{\pm}e^{-}$, $l^{\pm}$-hadron, and $\bar{p}p$ scattering; precision Z pole physics; and associated measurements of the W and top mass. They have verified the SM at the level of radiative corrections and strongly constrained the possibilities for new physics below the TeV scale (Yao *et al.*, 2006). There have been a number of global analyses of the constraints from precision electroweak on a possible Z' (Durkin and Langacker, 1986; London and Rosner, 1986; Amaldi *et al.*, 1987; del Aguila, Blair, *et al.*, 1987; Costa *et al.*, 1988; Gonzalez-Garcia and Valle, 1991; del Aguila *et al.*, 1992; Langacker and Luo, 1992; Langacker *et al.*, 1992; Cho *et al.*, 1998; Erler and Langacker, 1999, 2000; Chivukula and Simmons, 2002; Cacciapaglia *et al.*, 2006). Because of the number of different chiral fermions involved, it is difficult to do this in a model independent way, so most studies have focussed on specific classes of models, such as described in Sec. III.A, and have emphasized electroweak scale couplings and family universal charges.

Low-energy WNC experiments are affected by Z' exchange, which is mainly sensitive to its mass, and by Z-Z' mixing. Prior to the Tevatron and LEP 2 they yielded the best limits on the Z' mass. The Z-pole experiments at LEP and Stanford Linear Collider (SLC), on the other hand, are mainly sensitive to Z-Z' mixing, which lowers the mass of the Z relative to the SM prediction, and also modifies the $Z\bar{f}f$ vertices.

The effective four-Fermi Lagrangian for the WNC obtained from Eq. (9) is

$$-L_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \sum_{\alpha=1}^{n+1} \rho_\alpha \left[\sum_{\beta=1}^{n+1} \frac{g_\beta}{g_1} U_{\alpha\beta} J_\beta^\mu \right]^2, \tag{59}$$

where $\rho_\alpha \equiv M_W^2/M_\alpha^2 \cos^2\theta_W$, M_α are the mass eigenvalues, U is the orthogonal transformation defined in Eq. (17), and the currents are given in Eq. (10) (kinetic mixing can be added). Specializing to the $n=1$ case, this is

$$-L_{\text{eff}} = \frac{4G_F}{\sqrt{2}} (\rho_{\text{eff}} J_1^2 + 2wJ_1J_2 + yJ_2^2), \tag{60}$$

in the notation of Durkin and Langacker (1986) and Langacker and Luo (1992), where

$$\rho_{\text{eff}} = \rho_1 \cos^2\theta + \rho_2 \sin^2\theta,$$

$$w = \frac{g_2}{g_1} \cos\theta \sin\theta (\rho_1 - \rho_2),$$

$$y = \left(\frac{g_2}{g_1}\right)^2 (\rho_1 \sin^2\theta + \rho_2 \cos^2\theta), \tag{61}$$

with the mixing angle θ defined in Eq. (22). For small ρ_2 and θ, these are given by

$$\rho_{\text{eff}} \sim \rho_1, \quad w \sim \hat{\theta}, \quad y \sim \hat{\rho}_2, \tag{62}$$

where

$$\hat{\theta} \equiv \frac{g_2}{g_1}\theta = C\hat{\rho}_2, \quad \hat{\rho}_2 \equiv \left(\frac{g_2}{g_1}\right)^2 \rho_2. \tag{63}$$

C is the Higgs-dependent mixing parameter of $O(1)$ defined in Eq. (26). In the same limit, from Eq. (24),

$$\rho_1 \sim \rho_0(1 + \rho_0\theta^2/\rho_2) \xrightarrow[\rho_0=1]{} 1 + \theta^2/\rho_2 = 1 + C^2\hat{\rho}_2, \tag{64}$$

where ρ_0, defined in Eq. (16), is 1 if there are only Higgs singlets and doublets.

At the Z pole, in addition to the shift in M_1 below the SM value, any mixing will affect the current $\Sigma_\beta g_\beta U_{1\beta} J_\beta^\mu / g_1$ that couples to the Z_1. For $n=1$, the vector and axial couplings V_i and A_i of the Z_1 to fermion f_i, which determine the various Z pole asymmetries and partial widths (Yao *et al.*, 2006), become

$$V_i = \cos\theta g_V^1(i) + \frac{g_2}{g_1} \sin\theta g_V^2(i) \sim g_V^1(i) + \hat{\theta} g_V^2(i),$$

$$A_i = \cos\theta g_A^1(i) + \frac{g_2}{g_1} \sin\theta g_A^2(i) \sim g_A^1(i) + \hat{\theta} g_A^2(i), \tag{65}$$

where $g_{V,A}^\alpha(i)$ are defined in Eq. (10). It should be noted that the S,T,U formalism (Peskin and Takeuchi, 1990) only describes propagator corrections and is not appropriate for most Z''s.

2. Radiative corrections

The expressions for the electroweak couplings in Eqs. (10), (61), and (65) and for M_{Z_0} in Eq. (15) are valid at tree level only. One must also apply full radiative corrections. In practice, since one is searching for very small tree-level effects from the Z' it is a reasonable approximation to use the SM radiative corrections (Yao *et al.*, 2006) and neglect the effects of the Z' in loops.[5] However, some care is necessary in the definitions of the

[5]The largest effects are from Z_2 loops in μ decay, which modifiy slightly the relation between the extracted Fermi constant and the W and Z masses (Degrassi and Sirlin, 1989). Z_2 loops can also modify the relation between μ and β decay and therefore affect the Cabibbo-Kobayashi-Moskawa universality tests (Marciano and Sirlin, 1987). However, these effects are small for the currently allowed masses.

TABLE IV. 95% C.L. lower limits on various extra Z' gauge boson masses (GeV) and 90% C.L. ranges for the mixing $\sin\theta$ from precision electroweak data (columns 2–4), Tevatron searches (assuming decays into SM particles only), and LEP 2. The Tevatron numbers in parentheses are preliminary CDF results from March, 2008 based on 2.5 fb^{-1} (CDF note CDF/PUB/EXOTIC/PUBLIC/9160). From Erler and Langacker, 1999; Alcaraz *et al.*, 2006; Yao *et al.*, 2006; Aaltonen *et al.*, 2007.

	ρ_0 free	$\rho_0=1$	$\sin\theta(\rho_0=1)$	Tevatron	LEP 2
χ	551	545	(−0.0020) – (+0.0015)	822 (864)	673
ψ	151	146	(−0.0013) – (+0.0024)	822 (853)	481
η	379	365	(−0.0062) – (+0.0011)	891 (933)	434
LR	570	564	(−0.0009) – (+0.0017)	630	804
Sequential	822	809	(−0.0041) – (+0.0003)	923 (966)	1787

renormalized parameters, e.g., by using the modified minimal subtraction scheme rather than the on-shell definition of $\sin^2\theta_W$, to ensure that they are not significantly affected by Z' effects (Degrassi and Sirlin, 1989; Chankowski *et al.*, 2006).

3. Results

The results from precision electroweak and other data are shown for some specific models in Table IV and Fig. 1. One sees that the precision data strongly constrain the Z-Z' mixing angle θ. They also give lower limits on M_2, but these are weaker than the Tevatron and LEP 2 limits. The precision limit on the Z_ψ mass is low due to its weak coupling to the neutrino and its purely axial coupling to the e^-. There is no significant indication for a Z' in the data [although the NuTeV anomaly could possibly be explained by a Z' coupling to B-3L (Davidson *et al.*, 2002)]. The precision results are presented for two cases: ρ_0 free is for an arbitrary Higgs structure, while $\rho_0=1$ is for Higgs doublets and singlets with unrestricted charges (i.e., C is left free). There is little difference between the limits obtained. The precision electroweak constraints are for the g_2 value in Eq. (44) (except for the sequential model, which uses $g_2=g_1\sim 0.74$); for other values the limits on θ and M_2 scale as g_2^{-1} and g_2, respectively.

The stringent mixing limits from (mainly) the Z pole data lead to strong indirect limits on the Z' mass for specific theoretical values of C, as can be seen from the theoretical curves labeled 0, 1, 5, ∞ in Fig. 1 (Langacker and Luo, 1992). For the χ and LR models the label refers to the value of $|x|^2/(|\nu_u|^2+|\nu_d|^2)$, where $x/\sqrt{2}$ is the VEV of an extra Higgs doublet that is sometimes considered (transforming like an L doublet for χ or like the δ_L^0 defined in Sec. III.A.2 for the LR). The most commonly studied cases are for $x=0$, which yield $M_{Z_\chi} > 1368$ GeV, $M_{Z_{LR}} > 1673$ GeV at 95% C.L. For the ψ and η models, the label represents $\tan^2\beta \equiv |\nu_u|^2/|\nu_d|^2$, with $x=0$ assumed.

B. Constraints from colliders

1. Hadron colliders

The primary discovery mode for a Z' at a hadron collider is the Drell-Yan production of a dilepton resonance $pp(\bar{p}p)\to Z'\to \ell^+\ell^-$, where $\ell=e$ or μ (Langacker *et al.*, 1984; Barger *et al.*, 1987; del Aguila *et al.*, 1989; Dittmar, 1997; Leike, 1999; Godfrey, 2002; Carena *et al.*, 2004; Dittmar *et al.*, 2004; Kang and Langacker, 2005; Weiglein *et al.*, 2006; Yao *et al.*, 2006, Aaltonen *et al.*, 2007). Other channels, such as $Z'\to jj$ where j=jet (Weiglein *et al.*, 2006), $\bar{t}t$ (Han, Valencio, and Wung, 2004), $e\mu$ (Abulencia *et al.*, 2006), or $\tau^+\tau^-$, are also possible. The forward-backward asymmetry for $pp(\bar{p}p)\to\ell^+\ell^-$ (as a function of rapidity y for pp) due to γ, Z, Z' interference below the Z' peak is also important (Langacker *et al.*, 1984; Rosner, 1996a; Dittmar, 1997; Aaltonen *et al.*, 2007).

The cross section for hadrons A and B at center of mass energy $\sqrt{s}$ to produce a Z_α of mass M_α at rapidity y is, in the narrow width approximation (Langacker *et al.*, 1984),

$$\frac{d\sigma}{dy} = \frac{4\pi^2 x_1 x_2}{3M_\alpha^3}\sum_i [f_{q_i}^A(x_1)f_{\bar{q}_i}^B(x_2) + f_{\bar{q}_i}^A(x_1)f_{q_i}^B(x_2)] \times\Gamma(Z_\alpha\to q_i\bar{q}_i), \quad (66)$$

where $f_{q_i,\bar{q}_i}^{A,B}$ are the structure functions of quark (or antiquark) q_i ($\bar{q}_i$) in hadrons A or B, and the momentum fractions are

$$x_{1,2} = (M_\alpha/\sqrt{s})e^{\pm y}. \quad (67)$$

Neglecting mixing effects the decay width into fermion f_i is

$$\Gamma_{f_i}^\alpha \equiv \Gamma(Z_\alpha\to f_i\bar{f}_i) = \frac{g_\alpha^2 C_{f_i} M_\alpha}{24\pi}[\epsilon_L^\alpha(i)^2 + \epsilon_R^\alpha(i)^2], \quad (68)$$

where the fermion mass has been neglected. C_{f_i} is the color factor (1 for color singlets, 3 for triplets). Formulas including fermion mass effects, decays into bosons, Majorana fermions, etc., are given in Kang and Langacker (2005).

To a good first approximation, Eq. (66) leads to the Z' total production cross section (Leike, 1999)

$$\sigma_{Z'} = \frac{1}{s}c_{Z'}CK\exp\left(-A\frac{M_{Z'}}{\sqrt{s}}\right), \quad (69)$$

where $C=600$ (300) and $A=32$ (20) for pp ($p\bar{p}$) collisions, and $K\sim 1.3$ is from higher-order corrections. From Eq.

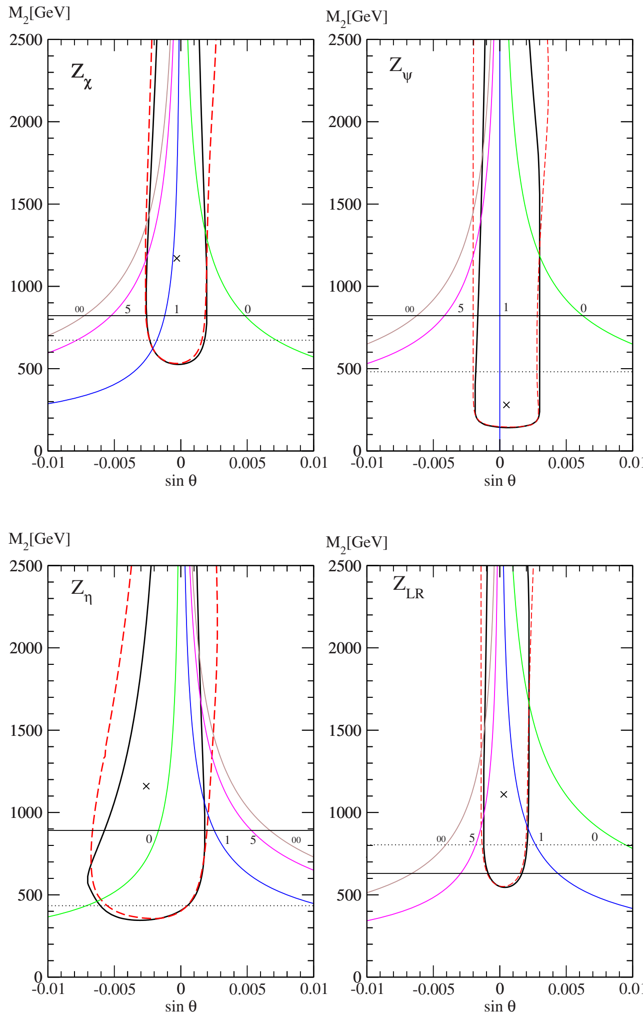

FIG. 1. (Color online) Limits on the Z' mass M_2 and the Z-Z' mixing angle θ for the χ, ψ, η, and LR (α=1.53) models. The solid (dashed) contours are 90% C.L. exclusions from precision electroweak data for ρ_0=1 (ρ_0=free). A cross $\times$ is the best fit. The horizontal solid line is the 95% C.L. Tevatron lower limit, assuming decays into SM particles only. The horizontal dotted line is the 95% C.L. lower limit from LEP 2. The contours marked 0, 1, 5, ∞ are for various theoretical relations between the mass and mixing and are defined in the text. Updated from Erler and Langacker, 1999.

(69), the predicted cross section falls exponentially as a function of $M_{Z'}$. The details of the Z' model are collected in $c_{Z'}$, which depends on $M_{Z'}$, the Z' couplings, and the masses of the decay products,

$$c_{Z'} \equiv \frac{4\pi^2}{3}\frac{\Gamma_{Z'}}{M_{Z'}}\left(B_u + \frac{1}{C_{ud}}B_d\right), \tag{70}$$

where C_{ud}=2, (25) $\Gamma_{Z'}$ is the total Z' width, and $B_f = \Gamma_f/\Gamma_{Z'}$ is the branching ratio into $f\bar{f}$. It is also useful to define

$$\sigma^f_{Z'} \equiv \sigma_{Z'}B_f = N_f/\mathcal{L}, \tag{71}$$

where N_f is the number of produced $f\bar{f}$ pairs for integrated luminosity $\mathcal{L}$. More detailed estimates for the Tevatron and LHC are given in Leike (1999), Godfrey (2002), Carena *et al.* (2004), Dittmar *et al.* (2004), and Fuks *et al.* (2008), including discussions of parton distribution functions, higher-order effects, width effects, resolutions, and backgrounds.

The production cross sections, widths, and branching ratios are considered in detail in Langacker *et al.* (1984), Barger *et al.* (1987), Gherghetta *et al.* (1998), and Kang and Langacker (2005). For the E_6 models, the total width is close to $0.01M_{Z'}$ assuming decays into SM fermions only and $g_2 \sim \sqrt{5/3}g \tan\theta_W$. However, $\Gamma_{Z'}$ would be larger if superpartners and/or exotics are light enough to be produced in the Z' decays, and could therefore be as large as $0.05M_{Z'}$ in the E_6 models (Kang and Langacker, 2005). The rates for a given channel, such as $\sigma^e_{Z'}$, decrease as $\Gamma^{-1}_{Z'}$ in that case. On the other hand, for smaller $\Gamma_{Z'}$ but fixed branching ratios (e.g., from some of the decoupled models described in Sec. III.D.1) the leptonic rate would decrease and the peak could be smeared out by detector resolution effects.

The Tevatron limits from the CDF and D0 Collaborations (Yao *et al.*, 2006; Aaltonen *et al.*, 2007) (dominated by the CDF e^+e^- search using 1.3 fb^{-1} of data) are given in Table IV. Figure 2 shows the sensitivity of the Tevatron and LHC to the E_6 bosons as a function of θ_{E_6} for $\mathcal{L}$=1 or 3 fb^{-1} (Tevatron), and 100 or 300 fb^{-1} (LHC), requiring 10 events in the combined e^+e^- and $\mu^+\mu^-$ channels. The Tevatron sensitivity is in the 600–900 GeV range for decays into standard model fermions only, but lower by as much as 200 GeV in the (extreme) case of unsuppressed decays into sparticles and exotics. The LHC sensitivity is around 4–5 TeV, but can be lower by ~1 TeV if the sparticle or exotic channels are open.

2. e^+e^- colliders

Z''s much heavier than the center of mass energy in e^+e^- collisions above the Z pole would manifest themselves as new four-fermion interactions analogous to Eq. (59), but with the α sum starting at 2. These would interfere with the virtual γ and Z contributions for leptonic and hadronic final states [see, e.g., Cheung (2001)].

The ALEPH, DELPHI, L3, and OPAL Collaborations at LEP2 have measured production cross sections and angular distributions or asymmetries for $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$, $\bar{c}c$, and $\bar{b}b$, as well as hadronic cross sections, at center of mass energies up to ~209 GeV (Alcaraz *et al.*, 2006). They saw no indication of new four-fermion interactions, and the combined lower limits for typical models are given in Table IV and Fig. 1.

Similarly, a future linear collider would have sensitivity to $M_{Z'}$ well above the center of mass energy by interference with the γ and Z (del Aguila and Cvetic, 1994; Leike and Riemann, 1997; Babich *et al.*, 1999; Richard, 2003; Godfrey *et al.*, 2005; Weiglein *et al.*, 2006). Observables could include production cross sections, forward-backward (FB) asymmetries, polarization (LR) asymmetries, and mixed FB-LR asymmetries for $e^+e^- \to e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$, $\bar{c}c$, $\bar{b}b$, and $\bar{t}t$; τ polarization; and

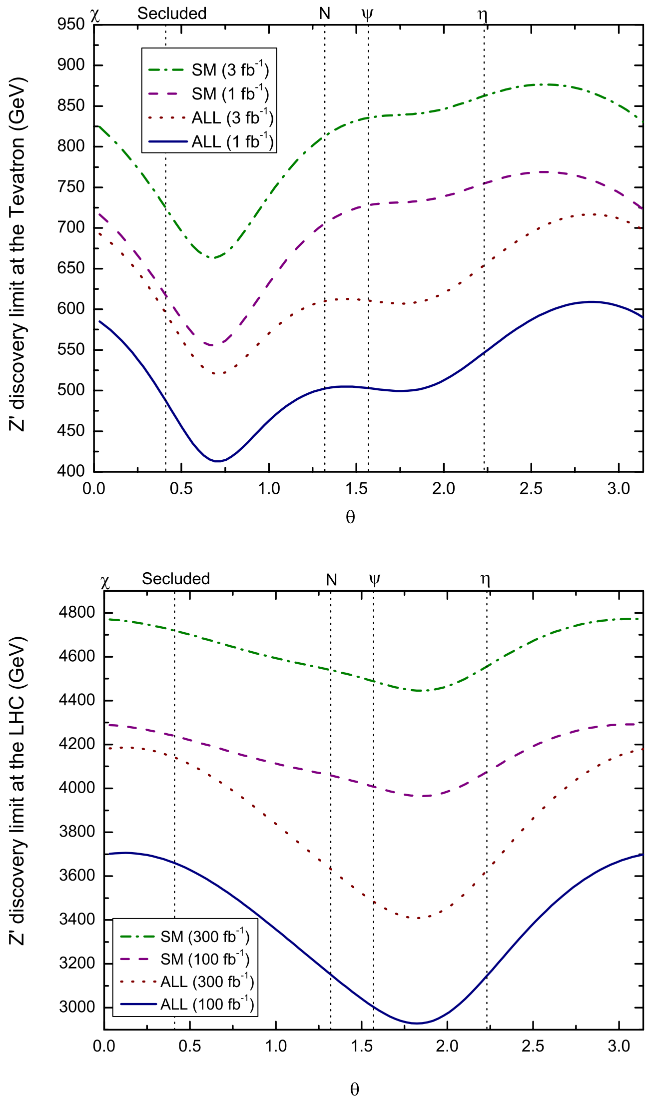

FIG. 2. (Color online) Discovery limits for an E_6 Z' as a function of $\theta \equiv \theta_{E_6}$ corresponding to a total of ten e^+e^- or $\mu^+\mu^-$ events using $\sigma_{Z'}$ from Eq. (69). In each panel the top two curves assume decays into SM fermions only, while the bottom two assume that decays into exotics and sparticles are unsuppressed. The different shapes of the Tevatron and LHC curves is because the u quark dominates at the Tevatron, while the u and d are more comparable at the LHC. From Kang and Langacker, 2005.

cross sections and polarization asymmetries for $\bar{q}q$. High luminosity, e^- polarization, and efficient tagging of heavy flavors are important. For example, the International Linear Collider with $\sqrt{s}=500$ GeV, $\mathcal{L}=1000$ fb^{-1}, and $P_{e^-}=80\%$ would have 5σ sensitivity to the E_6 and LR bosons in the range 2–5 TeV, increasing by $\sim$1 TeV for $\sqrt{s}=1$ TeV (Weiglein *et al.*, 2006). There is some chance that a Z' could be observed first at the International Linear Collider if its mass were beyond the LHC range or its couplings weak, in which case only M_2/g_2 could be determined for large M_2. More likely, the Z' would be discovered first and $M_{Z'}$ determined independently at the LHC or Tevatron. A GigaZ (Z-pole) option for the International Linear Collider would be extremely sensitive to Z-Z' mixing (Weiglein *et al.*, 2006).

C. Diagnostics of Z' couplings

Following the discovery of a resonance in the $\ell^+\ell^-$ channels, the next step would be to establish its spin-1 nature (as opposed to a spin-0 Higgs resonance or a spin-2 Kaluza-Klein graviton excitation). This can be done by the angular distribution in the resonance rest frame, which for spin 1 is

$$\frac{d\sigma^f_{Z'}}{d\cos\theta^*} \propto \frac{3}{8}(1+\cos^2\theta^*) + A^f_{FB}\cos\theta^*, \tag{72}$$

where θ^* is the angle between the incident quark or lepton and fermion f. Of course, for a hadron collider one does not know which hadron is the source of the q and which the $\bar{q}$ on an event by event basis, but the ambiguity washes out in the determination of the $1+\cos^2\theta^*$ distribution characteristic of spin 1 (Langacker *et al.*, 1984; Dittmar, 1997). The spin can also be probed in e^+e^- by polarization asymmetries (Weiglein *et al.*, 2006).

One would next want to determine the chiral couplings to the quarks, leptons, and other particles in order to discriminate between models. [The gauge coupling g_2 can be fixed to the value in Eq. (44), or alternatively can be taken as a free parameter if the charges are normalized by some convention.] This should be possible for masses up to $\sim$2–2.5 TeV at the LHC assuming typical couplings, but for higher masses there are too few events for meaningful diagnostics.

In the main LHC production channels, $pp \to Z' \to \ell^+\ell^-$ ($\ell=e,\mu$), one would be able to measure the mass $M_{Z'}$, the width $\Gamma_{Z'}$, and the leptonic cross section $\sigma^\ell_{Z'} = \sigma_{Z'}B_\ell$. By itself, $\sigma^\ell_{Z'}$ is not a useful diagnostic for the Z' couplings to quarks and leptons: while $\sigma_{Z'}$ can be calculated to within a few percent for given Z' couplings, the branching ratio into leptons B_ℓ depends strongly on the contribution of exotics and sparticles to $\Gamma_{Z'}$ (Kang and Langacker, 2005). However, $\sigma^\ell_{Z'}$ would be a useful indirect probe for the existence of the exotics or superpartners. Furthermore, the product $\sigma^\ell_{Z'}\Gamma_{Z'} = \sigma_{Z'}\Gamma_\ell$ does probe the absolute magnitude of the quark and lepton couplings.

The most useful diagnostics involve the relative strengths of Z' couplings to ordinary quarks and leptons. The forward-backward asymmetry as a function of the Z' rapidity $A^f_{FB}(y)$ (Langacker *et al.*, 1984) avoids the $\bar{q}q$ ambiguity in Eq. (72). For $AB \to Z' \to \bar{f}f$, define θ_{CM} as the angle of fermion f with respect to the direction of hadron A in the Z' rest frame, and let F (B) be the cross section for fixed rapidity y with $\cos\theta_{CM} > 0$ (<0). Then, $A^f_{FB}(y) \equiv (F-B)/(F+B)$, with

$$F \pm B \sim \begin{bmatrix} 4/3 \\ 1 \end{bmatrix} \sum_i [f^A_{q_i}(x_1)f^B_{\bar{q}_i}(x_2) \pm f^A_{\bar{q}_i}(x_1)f^B_{q_i}(x_2)] \times [\epsilon_L(q_i)^2 \pm \epsilon_R(q_i)^2][\epsilon_L(f)^2 \pm \epsilon_R(f)^2]. \tag{73}$$

Clearly, $A_{FB}^{f}(y)$ vanishes for pp at $y=0$, but can be nonzero at large y where there is more likely a valence q from the first proton and sea $\bar{q}$ from the other. The leptonic forward-backward asymmetry is sensitive to a combination of quark and lepton chiral couplings and is a powerful discriminant between models (Langacker *et al.*, 1984).

There are a number of additional probes. The ratio of cross sections in different rapidity bins (del Aguila *et al.*, 1993) gives information on the relative u and d couplings. Possible observables in other two-fermion final state channels include the polarization of produced τ's (Anderson *et al.*, 1992) and the $pp\to Z'\to jj$ cross section (Rizzo, 1993a; Weiglein *et al.*, 2006). There are no current plans for polarization at the LHC, but polarization asymmetries at a future or upgraded hadron collider would provide another useful diagnostic (Fiandrino and Taxil, 1992).

In four-fermion final state channels the rare decays $Z'\to Vf_1\bar{f}_2$, where $V=W$ or Z is radiated from the Z' decay products, have a double logarithmic enhancement. In particular, $Z'\to W\ell\nu_\ell$ (with $W\to$hadrons and an $\ell\nu_\ell$ transverse mass >90 GeV to separate from SM background) may be observable and projects out the left-chiral lepton couplings (Rizzo, 1987; Cvetic and Langacker, 1992a; Hewett and Rizzo, 1993). Similarly, the associated productions $pp\to Z'V$ with $V=(Z,W)$ (Cvetic and Langacker, 1992b) and $V=\gamma$ (Rizzo, 1993b) could yield information on the quark chiral couplings.

Finally, decays into two bosons, such as $Z'\to W^+W^-, Zh$, or $W^\pm H^\mp$, can occur only by Z-Z' mixing or with amplitudes related to the mixing. However, this suppression may be compensated for the longitudinal modes of the W or Z by the large polarization vectors, with components scaling as $M_{Z'}/M_W$ (Barger and Whisnant, 1987; del Aguila, Quiros, and Zwirner, 1987; Deshpande and Trampetic, 1988). For example, $\Gamma(Z'\to W^+W^-)\sim\theta^2$, which appears to be hopelessly small to observe. However, the enhancement factor is $\sim(M_{Z'}/M_W)^4$. Thus, from Eq. (26) these factors compensate, leaving a possibly observable rate that in principle could give information on the Higgs charges. In the limit $M_{Z'}\gg M_Z$ one has

$$\Gamma(Z'\to W^+W^-)=\frac{g_1^2\theta^2M_{Z'}}{192\pi}\left(\frac{M_{Z'}}{M_Z}\right)^4=\frac{g_2^2C^2M_{Z'}}{192\pi}. \quad (74)$$

Global studies of the possible LHC diagnostic possibilities for determining ratios of chiral charges in a model independent way and discriminating models are given in del Aguila *et al.* (1993) and Cvetic and Godfrey (1995). The complementarity of LHC and ILC observations is emphasized in del Aguila and Cvetic (1994), del Aguila, Cvetic, and Langacker (1995), Cvetic and Godfrey (1995), and Weiglein *et al.* (2006).

V. IMPLICATIONS

A. μ problem and extended Higgs/neutralino sectors

1. μ problem

As described in Sec. II.E.3, the μ problem of the MSSM can be solved in *singlet extended* models in which a symmetry forbids an elementary μ term, but allows a dynamical $\mu_{\text{eff}}=\lambda_S\langle S\rangle$. There are a number of realizations of this mechanism [see Accomando *et al.* (2006) and Barger *et al.* (2007c) for reviews]. The best known is the next to minimal model (NMSSM), in which a discrete Z_3 symmetry forbids μ but allows the cubic terms $\lambda_S SH_uH_d$ and $\kappa S^3/3$ in the superpotential (Ellis *et al.*, 1989). The original form of the NMSSM suffers from cosmological domain wall problems because of the discrete symmetry. This can be remedied in more sophisticated forms involving an R symmetry (Accomando *et al.*, 2006). A variation on that approach yields the new minimal model (nMSSM), in which the cubic term and its soft analog are replaced by tadpole terms linear in S with sufficiently small coefficients (Panagiotakopoulos and Tamvakis, 1999). A U(1)$'$ symmetry, which is perhaps more likely to emerge from a string construction, is another possibility (Suematsu and Yamagishi, 1995; Cvetic and Langacker, 1996a; Cvetic *et al.*, 1997). This avoids the domain wall problem by embedding the discrete symmetry of the NMSSM into a continuous one.

2. Extended Higgs sector

Conventional U(1)$'$ models necessarily involve extended Higgs sectors associated with the SM singlet fields whose VEVs break the U(1)$'$ symmetry. Especially interesting in this respect are those supersymmetric models involving a dynamical $\mu_{\text{eff}}=\lambda_S\langle S\rangle$. If one ignores Higgs sector CP violation,[6] then there will be an additional Higgs scalar associated with S that can mix with the two MSSM scalars from $H_{u,d}^0$. (There is also an additional pseudoscalar in the models involving a discrete symmetry.) Since the S does not couple directly to the SM fermions or gauge bosons, the LEP *lower* limits on the Higgs mass ($m_H>114.4$ GeV for the SM Higgs, and somewhat weaker in the MSSM) are weakened if the lightest Higgs has a significant singlet component (Barger *et al.*, 2006). Conversely, the theoretical *upper* limit on the lightest Higgs is also relaxed, from $\sim$130 GeV in the MSSM to around 170 GeV in the simplest U(1)$'$ model, due to the new F and D term contributions to the potential in Eq. (55). [One must include the loop corrections to these estimates (Barger *et al.*, 2006).] These relaxed limits allow lower values for $\tan\beta\equiv\nu_u/\nu_d$ in the U(1)$'$ models than are favored for the MSSM.

The UMSSM is the U(1)$'$ with a single S, with the potential in Eq. (55). In the decoupling limit, $\langle S\rangle\to\infty$

[6]Loop effects may generate significant CP effects, especially for the heavier Higgs states (Demir and Everett, 2004).

with $\mu_{\rm eff}$ fixed, the UMSSM reduces to the MSSM. Existing constraints favor this limit [unless the Z' is leptophobic, with small Z-Z' mixing due to a cancellation of the two terms in Δ^2 in Eq. (20)]. For large $\langle S\rangle$ the extra Higgs is heavy and mainly singlet (Barger *et al.*, 2006), so the Higgs sector is MSSM-like. However, more general U(1)′ models such as the secluded model in Sec. III.E.3, as well as other models such as the nMSSM, can yield significant doublet-singlet mixing, light singlet-dominated states, etc. (Erler *et al.*, 2002; Han, Langacker, and McElrath, 2004). [In fact, the secluded model reduces to the nMSSM in an appropriate limit (Barger *et al.*, 2006).] This may yield such nonstandard collider signatures as light weakly coupled Higgs, multiple Higgs with reduced couplings, and invisible decays into light neutralinos (Barger *et al.*, 2007a).

3. Extended neutralino sector

The neutralino sector of the MSSM (the $\tilde{B}$ and $\tilde{W}^0$ with soft masses $M_{\tilde{B},\tilde{W}}$; and two neutral Higgsinos $\tilde{H}^0_{u,d}$) is extended in U(1)′ models by one or more singlinos, $\tilde{S}$, and by the Z' gaugino $\tilde{Z}'$ with soft mass $M_{\tilde{Z}'}$ (Suematsu, 1998; Hesselbach *et al.*, 2002; Barger *et al.*, 2005, 2007b; Choi *et al.*, 2007). (There could also be soft mass or kinetic $\tilde{B}$-$\tilde{Z}'$ mixing terms.)

In the $(\tilde{B},\tilde{W}^0,\tilde{H}^0_d,\tilde{H}^0_u,\tilde{S},\tilde{Z})$ basis, the mass matrix for the six neutralinos in the UMSSM is

$$M_{\chi^0}=\begin{pmatrix} M_{\tilde{B}} & 0 & -g'\nu_d/2 & g'\nu_u/2 & 0 & 0\\ 0 & M_{\tilde{W}} & g\nu_d/2 & -g\nu_u/2 & 0 & 0\\ -g'\nu_d/2 & g\nu_d/2 & 0 & -\mu_{\rm eff} & -\mu_{\rm eff}\nu_u/s & g_2Q_d\nu_d\\ g'\nu_u/2 & -g\nu_u/2 & -\mu_{\rm eff} & 0 & -\mu_{\rm eff}\nu_d/s & g_2Q_u\nu_u\\ 0 & 0 & -\mu_{\rm eff}\nu_u/s & -\mu_{\rm eff}\nu_d/s & 0 & g_2Q_Ss\\ 0 & 0 & g_2Q_d\nu_d & g_2Q_u\nu_u & g_2Q_Ss & M_{\tilde{Z}'}\end{pmatrix}. \tag{75}$$

In the decoupling limit with $g_2Q_Ss \gg M_{\tilde{Z}'}$ the singlino and the Z' gaugino will combine to form an approximately Dirac fermion with mass $g_2Q_Ss\sim M_{Z'}$ and little mixing with the four MSSM neutralinos. For large $M_{\tilde{Z}'} \gg g_2Q_Ss$, on the other hand, there will be a heavy Majorana $\tilde{Z}'$, and a much lighter singlino with a seesaw type mass $\sim -M^2_{Z'}/M_{\tilde{Z}'}$. For smaller s there can be significant mixing with the MSSM neutralinos. One can easily extend to secluded models (Erler *et al.*, 2002; Han, Langacker, and McElrath, 2004), models with multiple U(1)′'s (Hesselbach *et al.*, 2002), or singlet extended models with discrete symmetries (Barger *et al.*, 2005, 2007c; Accomando *et al.*, 2006). In many of these cases there are light singlino-dominated states [which can be the lightest supersymmetric particle (LSP)] and/or significant mixing effects. These can lead to a variety of collider signatures very different from the MSSM (Barger *et al.*, 2007b). For example, in some cases there are four MSSM-like neutralinos with production and cascades similar to the MSSM. However, the lightest of these may then undergo an additional decay to a singlino LSP, accompanied by an on-shell Z or Higgs. Enhanced rates for the decay of chargino-neutralino pairs to three or more leptons are possible. It is also possible for the lightest Higgs to decay invisibly to two light singlinos. Cold dark matter implications are described in Sec. V.G.

B. Exotics

Almost all U(1)′ models require the addition of new chiral exotic states to cancel anomalies (Sec. II.C). Precision electroweak constraints favor that these are *quasichiral*, i.e., vector pairs under the SM but chiral under U(1)′. Examples are the SU(2)-singlet D, D^c quarks with charge $-1/3$ in the E_6 model (Table II); the SU(2) doublet pairs in E_6 which may be interpreted either as additional Higgs pairs $H_{u,d}$ or as exotic lepton doublets; or SM singlets. Realistic models must provide means of generating masses for such exotics, for example, by coupling to chiral (or nonchiral) singlets which acquire VEVs, such as SDD^c, and also for their decays.

Consider the example of the exotic D quarks, which can be pair produced by QCD processes at a hadron collider, and their scalar supersymmetric partners $\tilde{D}$, produced with an order of magnitude smaller cross section. (The rates are smaller for exotic leptons.) Once produced, there are three major decay possibilities (Kang *et al.*, 2008).

- The decay may be $D\to u_iW^-$, $D\to d_iZ$, or $D\to d_iH^0$, if driven by mixing with a light charge $-1/3$ quark (Barger *et al.*, 1986; Andre and Rosner, 2004). The current limit is $m_D \gtrsim 200$ GeV (Andre and Rosner, 2004), which should be improved to $\sim$1 TeV at the LHC. However, such mixing is forbidden in the supersymmetric E_6 model if R parity is conserved.

- One may have $\tilde{D}\to jj$ if there is a small diquark operator such as $u^c d^c D^c$, or $\tilde{D}\to j\ell$ for a leptoquark operator like LQD^c. (They cannot both be present because of proton decay.) Such operators do not by themselves violate R parity ($R=+1$ for the scalar), and therefore allow a stable lightest supersymmetric particle. They are strongly constrained by the K_L-K_S mass difference and by μ-e conversion, but may still be significant (Kang *et al.*, 2008). If the scalar $\tilde{D}$ is heavier than the fermion, then it may decay resonantly into the fermion pair, or into a D and neutralino (or gluino). The lighter fermion D can decay into a neutralino and nonresonant fermion pair via a virtual $\tilde{D}$ or via a real or virtual squark or slepton. A heavier fermion will usually decay into an on-shell $\tilde{D}$ and a neutralino (or gluino), with the $\tilde{D}$ decaying to fermions. The signals from these decays, especially for a heavier scalar, may be difficult to extract from normal supersymmetry cascades, especially for diquarks. However, there are some possibilities based on missing transverse energy, lepton multiplicities, p_T, etc. (Kang *et al.*, 2008).
- They may be stable at the renormalizable level due to the U(1)′, or to an accidental or other symmetry, so that they hadronize and escape from or stop in the detector (Kang *et al.*, 2008), with signatures (Kraan *et al.*, 2007) somewhat similar to the quasistable gluino expected in split supersymmetry (Arkani-Hamed and Dimopoulos, 2005). They could then decay by higher-dimensional operators on a time scale of $\lesssim 10^{-1}$–100 s, short enough to avoid cosmological problems (Kawasaki *et al.*, 2005). These operators could allow direct decays to SM particles, or they could involve SM singlets with VEVs which could induce small mixings with ordinary quarks.

Exotics carrying SM charges significantly modify the running of the SM couplings, and therefore can affect gauge unification unless they occur in SU(5)-type multiplets. Examples of U(1)′ constructions which preserve the MSSM running at tree level are described in Secs. III.A.3 and III.B.

C. The Z' as a factory

The decays of a Z' could serve as an efficient source of other particles if it is sufficiently massive. This has been explored in detail for slepton production, $pp\to Z' \to \tilde{\ell}\tilde{\ell}^*$, with $\tilde{\ell}\to\ell+\mathrm{LSP}$, assuming that $M_{Z'}$ is already known from the conventional $\ell^+\ell^-$ channel (Baumgart *et al.*, 2007). This can greatly extend the discovery reach of the $\tilde{\ell}$ and may give information on the identity of the LSP. Decays of the Z' could also be a useful production mechanism for pairs of exotics (Rosner, 1996b) or heavy Majorana neutrinos (Duncan and Langacker, 1986; del Aguila and Aguilar-Saavedra, 2007). The latter could lead to the interesting signature of like sign leptons +jets. The total width $\Gamma_{Z'}$, in combination with other constraints on the quark and lepton charges, would also give some information on the exotic or sparticle decays (Gherghetta *et al.*, 1998; Kang and Langacker, 2005).

D. Flavor changing neutral currents

In Sec. II it was implicitly assumed that the U(1)′ charges were family universal. This assumption implies that the Z' couplings are unaffected by fermion mixings and remain diagonal [the Glashow-Iliopoulos-Maiani (GIM) mechanism]. However, many models involve nonuniversal charges, as described in Sec. III.D.3. We rewrite the U(1)′ current in Eq. (10) as

$$J^\mu_\alpha = \bar{f}^0_L \gamma^\mu \epsilon^\alpha_{fL} f^0_L + \bar{f}^0_R \gamma^\mu \epsilon^\alpha_{fR} f^0_R, \tag{76}$$

where f^0_L is a column vector of weak-eigenstate left chiral fermions of a given type (i.e., u^0_L, d^0_L, e^0_L, or ν^0_L), and similarly for f^0_R. The ϵ^α_f are diagonal matrices of U(1)′ charges. The $f^0_{L,R}$ are related to the mass eigenstates $f_{L,R}$ by

$$f^0_L = V^{f\dagger}_L f_L, \quad f^0_R = V^{f\dagger}_R f_R, \tag{77}$$

where $V^f_{L,R}$ are unitary. In particular, the Cabibbo-Kobayashi-Maskawa (CKM) and Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrices are given by $V^u_L V^{d\dagger}_L$ and $V^\nu_L V^{e\dagger}_L$, respectively. In the mass basis,

$$J^\mu_\alpha = \bar{f}_L \gamma^\mu B^\alpha_{fL} f_L + \bar{f}_R \gamma^\mu B^\alpha_{fR} f_R, \tag{78}$$

where

$$B^\alpha_{fL} \equiv V^f_L \epsilon^\alpha_{fL} V^{f\dagger}_L, \quad B^\alpha_{fR} \equiv V^f_R \epsilon^\alpha_{fR} V^{f\dagger}_R. \tag{79}$$

For family universal charges, $\epsilon_{fL,R}$ are proportional to the identity, and $B_{fL,R}=\epsilon_{fL,R}$. However, for the nonuniversal case, $B_{fL,R}$ will in general be nondiagonal. As a simple two family example, if $\epsilon=\left(\begin{smallmatrix}0&0\\0&1\end{smallmatrix}\right)$ and V is a rotation of the same form as Eq. (22) then

$$J^\mu = \sin^2\theta \bar{f}_1\gamma^\mu f_1 + \cos^2\theta \bar{f}_2\gamma^\mu f_2 + \sin\theta\cos\theta(\bar{f}_1\gamma^\mu f_2 + \bar{f}_2\gamma^\mu f_1). \tag{80}$$

The formalism for FCNC mediated by Z', and also by off-diagonal Z couplings induced by Z-Z' mixing, was developed by Langacker and Plumacher (2000), and limits were obtained for a number of tree and loop level mixings and decays. The limits from K^0-$\bar{K}^0$ mixing (including CP violating effects) and from μ-e conversion in muonic atoms are sufficiently strong to exclude significant nonuniversal effects for the first two families for a TeV-scale Z' with electroweak couplings. However, nonuniversal couplings for the third family are still possible and could contribute (Langacker and Plumacher, 2000; Leroux and London, 2002; Barger, Chiang, Jiang, and Langacker, 2004; Barger *et al.*, 2004a, 2004b; Baek *et al.*, 2006; Chen and Hatanaka, 2006; Chiang *et al.*, 2006; He and Valencia, 2006; Cheung *et al.*, 2007) to processes such as $B\bar{B}$ and $D\bar{D}$ mixing, $B\to\mu^+\mu^-$, or $b\to ss\bar{s}$ (such as in $B \to \phi K$). Since the Z' effects are at tree level, they may be important even for small couplings since they are competing with SM or MSSM loop effects. The possible anomaly observed in the $Z \to \bar{b}b$ forward-backward asymmetry (Yao *et al.*, 2006) could possibly be a (flavor diagonal) result of a nonuniversal Z' coupling (Erler and Langacker, 2000). Collider processes such as single top production could possibly be observable as well (Arhrib *et al.*, 2006).

The nonuniversal couplings could also be relevant to loop effects, such as $b \to s\gamma$ or $\mu \to e\gamma$, or intrinsic magnetic or electric dipole moments. One interpretation of the possible anomaly (Yao *et al.*, 2006) suggested by the BNL experiment for the anomalous magnetic moment of the μ involves the vertex diagram with a Z' exchange [see, e.g., Cheung *et al.* (2007)]. The flavor-diagonal diagram with an internal μ is too small to be relevant (unless $M_{Z'} \sim 100$ GeV or the couplings are large). However, an internal τ enhances the effect by m_τ/m_μ, and the anomaly could be accounted for by a TeV-scale Z' with large μ-τ mixing.

Mixing between the ordinary and exotic fermions can also lead to FCNC effects (Langacker and London, 1988). For example, a small d^c-D^c mixing in the E_6 model of Table II would induce off diagonal couplings of the Z_2^0 to the light and heavy mass eigenstate, while a d-D mixing [i.e., between an SU(2) doublet and singlet] would generate similar effects for the ordinary Z^0. Off-diagonal vertices between the light mass eigenstates, such as $Z_\alpha^0 \bar{b} s$, would be induced as second-order effects.

E. Supersymmetry breaking, Z' mediation, and the hidden sector

U(1)$'$'s have many possible implications for supersymmetry breaking and mediation, and for communication with a hidden sector. For example, one limit of the single S scenario of Sec. III.E.2 requires large (TeV-scale) soft masses in the Higgs sector, suggesting the possibility of heavy sparticles as well (Everett *et al.*, 2000).

Another implication is the U(1)$'$ D term contribution to the scalar potential (Kolda and Martin, 1996),

$$V_D = \frac{1}{2} D^2 \equiv \frac{1}{2}\left(- g_2 \sum_i Q_i |\phi_i|^2 \right)^2. \tag{81}$$

V_D of course contributes to the minimization conditions and Higgs sector masses. Assuming a value $D^{\min} \neq 0$ for D at the minimum, it gives a contribution to the masses m_i^2 of the squarks, sleptons, and exotic scalars

$$\Delta m_i^2 = (-D^{\min})(g_2 Q_i). \tag{82}$$

For a single S field, one has $-D^{\min} = g_2(Q_u|\nu_u|^2 + Q_d|\nu_d|^2 + Q_S|s|^2)/2$ in the notation of Sec. II.B. Δm_i^2 can be of either sign, and must be added to other supersymmetric and soft contributions. When the U(1)$'$ scale is large, there is a danger of overall negative mass squares which destabilize the vacuum. However, in that case there is the possibility of breaking along a D flat direction in which $V_D^{\min}$ is small, as in the secluded models (Erler *et al.*, 2002). Positive D term contributions to the slepton masses have been suggested as a means of compensating the negative ones from anomaly mediated supersymmetry breaking (Murakami and Wells, 2003; Anoka *et al.*, 2004). The D term quartic interactions also contribute to the RGE equations for the soft masses [see, e.g., Cvetic *et al.* (1997) and Langacker and Wang (1998)].

U(1)$'$'s have been invoked in many models of supersymmetry breaking or mediation. For example, many models of gauge mediation involve a U(1)$'$ which may help transmit the breaking by loop effects and/or D terms in the hidden or ordinary sectors (Dobrescu, 1997; Cheng *et al.*, 1998, 1999; Kaplan *et al.*, 1999; Langacker *et al.*, 1999). The Fayet-Iliopoulos terms (Sec. III.F.3) associated with anomalous U(1)$'$'s in string constructions may also help trigger and transmit supersymmetry breaking (Mohapatra and Riotto, 1997; Dvali and Pomarol, 1998).

In many string constructions particles in both the ordinary and hidden sector may carry U(1)$'$ charges, allowing for the possibility of Z' *mediation* (Langacker *et al.*, 2008). The simplest case is that the U(1)$'$ gauge symmetry is not broken in the hidden sector, but the Z' gaugino acquires a mass from the supersymmetry breaking. The Z'-$\tilde{Z}'$ mass difference induces ordinary sector scalar masses at one loop, and SM gaugino masses at two loops. Requiring the latter to be in the range 10^2–10^3 GeV implies $M_{\tilde{Z}'} \gtrsim 10^3$ TeV (for electroweak couplings), with the sparticles, exotics, and Z' around 10–100 TeV and the electroweak scale obtained by a fine tuning, i.e., a version of split supersymmetry (Arkani-Hamed and Dimopoulos, 2005). String embeddings of this scenario are addressed in Verlinde *et al.* (2008). It can also be combined with other mediation scenarios, allowing a lower Z' scale (Nakayama, 2008). A Z' communicating with a hidden sector could also allow the production and decays into SM particles of relatively light *hidden valley* particles (Han *et al.*, 2007).

F. Neutrino mass

The seesaw model [see, e.g., Mohapatra *et al.* (2007)] leads to a small Majorana mass $m_\nu \sim -m_D^2/M_{\nu^c}$ for the ordinary doublet neutrinos ν, where m_D is a Dirac mass (generated by the VEV of a Higgs doublet), and $M_{\nu^c} \gg m_D$ is the Majorana mass of the heavy singlet ν^c,

$$-L_\nu = m_D \bar{\nu}_L \nu_R + \tfrac{1}{2} M_{\nu^c} \bar{\nu}_L^c \nu_R + \text{H.c.}, \tag{83}$$

where ν_R is the conjugate of ν_L^c. For $m_D \sim 100$ GeV and $M_{\nu^c} \sim 10^{14}$ GeV one obtains $|m_\nu|$ in the observed 0.1 eV range. However, if ν^c is charged under a U(1)$'$ then M_{ν^c} cannot be much larger than the U(1)$'$ scale. One possibility for a TeV-scale Z' is that ν^c is neutral, as in the N model (Ma, 1996; Barger *et al.*, 2003; Kang, Langacker, and Li, 2005; King *et al.*, 2006). Then, a conventional seesaw (Keith and Ma, 1996; Ma, 1996; Kang, Langacker, and Li, 2005; King *et al.*, 2006) and leptogenesis

(Hambye *et al.*, 2001) scenario can be possible if a large M_{ν^c} can be generated. For other models with TeV-scale $M_{Z'}$ one must invoke an alternative to the seesaw. For example, small Majorana masses can be generated using the double seesaw mechanism (involving an additional power of $M_{\nu^c}^{-1}$), or by invoking a Higgs triplet (Kang, Langacker, and Li, 2005; Mohapatra *et al.*, 2007).

Another possibility, which can lead to either small Dirac or Majorana masses, involves higher-dimensional operators (Cleaver *et al.*, 1998; Langacker, 1998; Arkani-Hamed, Hall, *et al.*, 2001; Borzumati and Nomura, 2001; Gogoladze and Perez-Lorenzana, 2002; Kang, Langacker, and Li, 2005; Chen *et al.*, 2007; Demir *et al.*, 2008). For example, a superpotential operator $W = SLH_u\nu^c/M$ could generate a small Dirac mass in the correct range for $\langle S\rangle \sim 10^6$ GeV and $M \sim 10^{18}$ GeV. Such a VEV can easily occur in intermediate scale models (Cleaver *et al.*, 1998; Langacker, 1998) or in the Z' mediation scenario (Langacker *et al.*, 2008). Higher powers could occur for a larger $\langle S\rangle$ or smaller M, associated, e.g., with an anomalous U(1)$'$ (Gogoladze and Perez-Lorenzana, 2002). Nonholomorphic (wrong Higgs) terms (see Sec. II.E.2) can also lead to naturally small Dirac masses, suppressed by the ratio of the supersymmetry breaking and mediation scales (Demir *et al.*, 2008). In some cases, a Z' gaugino is needed to generate a fermion mass at loop level from a nonholomorphic soft term. In these mechanisms, some low energy symmetry such as a U(1)$'$ must forbid a renormalizable level Dirac mass term $W = LH_u\nu^c$, while allowing the higher-dimensional operators. (The renormalizable level term is allowed in the E_6 models.) Discrete gauge symmetries (i.e., remnants of a gauge symmetry broken at a high scale) may also help restrict the allowed operators (Luhn and Thormeier, 2008).

Some mechanisms (Ma, 1996; Langacker, 1998) also allow the generation of light sterile neutrino masses and ordinary-sterile mixing, as suggested by the LSND experiment.

Recently, it was shown (Nelson and Walsh, 2008) that the LSND and MiniBooNE results could be reconciled in a model involving ordinary and sterile neutrinos if there is a very light (~ 10 keV) Z' coupled to B-L with a very weak coupling ($\lesssim 10^{-5}$). In analogy with the Mikheyev-Smirnov-Wolfenstein (MSW) effect (Mohapatra *et al.*, 2007) the Z' generates a potential in matter that is different for the ordinary and sterile neutrinos and strongly energy dependent.

The right-handed components of light Dirac neutrinos could upset the successful predictions of big bang nucleosynthesis if they were present in sufficient numbers. Mass and Yukawa coupling effects are too small to be dangerous. However, couplings of ν^c to a TeV-scale Z' could have kept them in equilibrium until relatively late (Olive *et al.*, 1981). A detailed estimate (Barger *et al.*, 2003) found that too much ^{4}He would have been produced for light Dirac neutrinos for most of the E_6 models unless $M_{Z'} \gtrsim 1$–3 TeV. Similar constraints follow from supernova cooling (Rizzo, 1991). These limits disappear, however, for couplings close to the N model. This is especially relevant for a parameter range of the generalized E_6 model [with two U(1)$'$'s] in which the Z_N is much lighter than the orthogonal boson, but nevertheless no Majorana masses are allowed (Kang, Langacker, and Li, 2005).

G. Cosmology

1. Cold dark matter

U(1)$'$ models (de Carlos and Espinosa, 1997; Barger, Kao, *et al.*, 2004; Burger *et al.*, 2007d; Lee, Matchev, and Nasri, 2007; Nakamura and Suematsu, 2007; Belanger *et al.*, 2008; Hur *et al.*, 2008; Pospelov *et al.*, 2008), as well as other singlet extended models with a dynamical μ term (Menon *et al.*, 2004; Accomando *et al.*, 2006; Barger *et al.*, 2007c, 2007d), have many implications for cold dark matter. For example, the extended neutralino sector in Eq. (75) allows the possibility of a light singlino as the LSP (de Carlos and Espinosa, 1997; Barger, Kao, *et al.*, 2004; Menon *et al.*, 2004; Barger *et al.*, 2007d; Nakamura and Suematsu, 2007), with efficient annihilation into a light Z' or into the Z (via small admixtures with the Higgsinos). More generally, the LSP may contain admixtures of $\tilde{S}$ or $\tilde{Z}'$ with the MSSM neutralinos, or allow a modified MSSM composition for the LSP. The models also have enlarged Higgs sectors and different allowed ranges, extending the possible mechanisms for Higgs-mediated LSP annihilation. Most of the interesting cases should be observable in direct detection experiments (Barger *et al.*, 2007d).

There are other LSP candidates in U(1)$'$ models. For example, the scalar partners $\tilde{\nu}^c$ of the singlet neutrinos become viable thermal cold dark matter candidates due to the possibility of annihilation through the Z' (Lee, Matchev, and Nasri, 2007). Other possibilities include a neutral exotic particle or multiple stable particles (Hur *et al.*, 2008), a heavy Dirac neutrino (Belanger *et al.*, 2008), or a semisecluded weak sector (Pospelov *et al.*, 2008) coupled via a Z'.

2. Electroweak baryogenesis

The seesaw model of neutrino mass allows the possibility of explaining the observed baryon asymmetry by leptogenesis, i.e., the decays of the heavy Majorana neutrino generate a small lepton asymmetry, which is partially converted to a baryon asymmetry by the electroweak sphaleron process (Mohapatra *et al.*, 2007). As discussed in Sec. V.F, however, an additional U(1)$'$ symmetry often forbids the seesaw model. Some of the alternatives discussed there allow other forms of leptogenesis (Chun, 2005).

However, the U(1)$'$ (Kang, Langacker, *et al.*, 2005; Ham and Oh, 2007; Ham *et al.*, 2007) and other singlet extended models (Menon *et al.*, 2004; Barger *et al.*, 2007c; Profumo *et al.*, 2007) open up the possibility of a completely different mechanism, *electroweak baryogenesis*. In this scenario, the interactions of particles with the expanding bubble wall from a strongly first order electroweak phase transition lead to a *CP* asymmetry, which is then converted to a baryon asymmetry by sphaleron processes. However, the SM does not have a strong first order transition or sufficient *CP* violation; the MSSM has only a small parameter range involving a light stop for the transition, and there is tension between the *CP* violation needed and electric dipole moment constraints (Carena, Quiros, *et al.*, 2003). In the extended models, however, there is a tree-level cubic scalar interaction [the $\lambda_S A_S S H_u^0 H_d^0$ term in Eq. (55)], which can easily lead to the needed strong first order transition. There are also possible new sources of tree-level *CP* violation in the Higgs sector (Kang, Langacker, *et al.*, 2005; Ham and Oh, 2007), which can contribute to the baryon asymmetry but have negligible effect on electric dipole moments.

3. Cosmic strings

A broken global or gauge U(1)′ can lead to *cosmic strings*, which are allowed cosmologically for a wide range of parameters and which could have interesting implications for gravitational waves, dark matter, particle emission, and gravitational lensing. For a recent discussion, with emphasis on breaking a supersymmetric U(1)′ along an almost flat direction, see Cui *et al.* (2008).

VI. CONCLUSIONS AND OUTLOOK

A new U(1)′ gauge symmetry is one of the best motivated extensions of the standard model. For example, U(1)′'s occur frequently in superstring constructions. If there is supersymmetry at the TeV scale, then both the electroweak and Z' scales are usually set by the scale of soft supersymmetry, so it is natural to expect $M_{Z'}$ in the TeV range. [One exception is when the U(1)′ breaking occurs along an approximately flat direction, in which case a large breaking scale could be associated with fermion mass hierarchies generated by higher-dimensional operators.] Similarly, TeV-scale U(1)′s (or Kaluza-Klein excitations of the photon and Z) frequently occur in models of dynamical symmetry breaking, little Higgs models, and models with TeV^{-1}-scale extra dimensions. Other constructions, such as nonsupersymmetric grand unified theories larger than SU(5), also lead to extra U(1)′s, but in these cases there is no particular reason to expect breaking at the TeV scale (and breaking below the GUT may lead to rapid proton decay).

The observation of a Z' would have consequences far beyond the existence of a new gauge boson. Anomaly cancellation would imply the existence of new fermions. These could be right-handed neutrinos, but usually there are additional particles with exotic electroweak quantum numbers. There must also be at least one new SM singlet scalar whose VEV breaks the U(1)′ symmetry. This scalar could mix with the Higgs doublet(s) and significantly alter the collider phenomenology. The Z' couplings could be family nonuniversal, allowing new tree-level contributions to t, b, and τ decays.

In the supersymmetric case the U(1)′ could solve the μ problem by replacing μ by a dynamical variable linked to the U(1)′ breaking, and the allowed MSSM parameter range would be extended. The singlets and exotics would be parts of chiral supermultiplets, and there would be extended neutralino sectors associated with the new singlino and gaugino, modifying the collider physics and cold dark matter possibilities. Gauge unification could be maintained if the exotics fell into SU(5)-type multiplets. The U(1)′ symmetry would also constrain the possibilities for neutrino mass and might be related to proton stability and R-parity conservation. A Z' might also couple to a hidden sector and could play a role in supersymmetry breaking or mediation. Finally, a dynamical μ would allow a strong first-order electroweak phase transition and new sources of *CP* violation in the Higgs sector, making electroweak baryogenesis more likely than in the SM or the MSSM, with the ingredients observable in the laboratory.

There are large classes of Z' models, distinguished by the chiral charges of the quarks, leptons, and Higgs fields, as well as the Higgs and exotic spectrum, gauge coupling, Z' mass, and possible mass and kinetic mixing. In string constructions, U(1)′'s that do not descend through SO(10) or left-right symmetry can have seemingly random charges. There is no simple classification or parametrization that takes into account all of the possibilities. One (model-independent) approach, valid for family universality, is to take a conventional value for the new gauge coupling, and regard the charges of the left-handed quarks (Q_L), leptons (L_L), and antiparticles u_L^c, d_L^c, and e_L^+, as well as $M_{Z'}$, $\Gamma_{Z'}$, and the mixing angle θ as free parameters relevant to experimental searches. However, eight parameters are too many for most purposes, so one must resort to specific models or lower-dimensional parametrizations to illustrate the possibilities. A recommended set are those summarized in Tables I–III.

Table I lists the $U(1)_{3R}\times U(1)_{BL}$ model, which is a one-parameter (not counting the Z' mass and mixing) set of models based on various forms of SO(10) and left-right symmetry, and a two-parameter generalization motivated by more general embeddings or by kinetic mixing. It requires no exotics other than ν_L^c. However, the supersymmetric version requires nonchiral Higgs doublets and (probably) vector pairs of SM singlets, and does not solve the μ problem.

Table II lists E_6-motivated models. A whole class of interesting models involves one free parameter θ_{E_6} or a two parameter generalization with kinetic mixing (or a third parameter if the gauge coupling is varied). These models illustrate typical exotics, and (with the exception of the χ model) the supersymmetric version involves a dynamical μ term. However, supersymmetric gauge unification requires an additional vector pair of Higgs-type doublets.

The models in Table III are examples of supersymmetric models with a dynamical μ that are consistent with gauge unification without additional vector pairs.

Three parameters, including the gauge coupling, are relevant to the nonexotic sector.

If there is a Z' with typical electroweak scale couplings to the ordinary fermions, it should be readily observable at the LHC for masses up to $\sim$4–5 TeV, or at the Tevatron for masses up to $\sim$600–900 GeV. Significant diagnostic probes of the Z' couplings would be possible up to $\sim$2–2.5 TeV. A future International Linear Collider would extend the range somewhat, and would provide complementary diagnostics. Within the context of supersymmetry, the observation of a Z' could completely alter the paradigm of having the MSSM at the TeV scale, with a desert up to a scale of grand unification or heavy Majorana neutrino masses, and would suggest a whole range of new laboratory and cosmological consequences. In the nonsupersymmetric case, a Z' might be one of the first experimental manifestations of a new TeV scale sector of physics.

ACKNOWLEDGMENTS

I am extremely grateful to Vernon Barger, Mirjam Cvetič, Jens Erler, and all of my other collaborators on work related to this article, and Hye-Sung Lee for a careful reading of the manuscript. This work was supported by the Friends of the IAS and by NSF Grant No. PHT-0503584.

REFERENCES

Aaltonen, T., *et al.*, (CDF), 2007, Phys. Rev. Lett. **99**, 171802.

Abulencia, A., *et al.*, (CDF), 2006, Phys. Rev. Lett. **96**, 211802.

Accomando, E., *et al.*, 2006, e-print arXiv:hep-ph/0608079.

Agashe, K., A. Delgado, M. J. May, and R. Sundrum, 2003, J. High Energy Phys. **08**, 050.

Agashe, K., *et al.*, 2007, Phys. Rev. D **76**, 115015.

Alcaraz, J., *et al.*, 2006, e-print arXiv:hep-ex/0612034.

Amaldi, U., *et al.*, 1987, Phys. Rev. D **36**, 1385.

Anastasopoulos, P., M. Bianchi, E. Dudas, and E. Kiritsis, 2006, J. High Energy Phys. **11**, 057.

Anastasopoulos, P., T. P. T. Dijkstra, E. Kiritsis, and A. N. Schellekens, 2006, Nucl. Phys. B **759**, 83.

Anderson, J. D., M. H. Austern, and R. N. Cahn, 1992, Phys. Rev. D **46**, 290.

Andre, T. C., and J. L. Rosner, 2004, Phys. Rev. D **69**, 035009.

Andrianov, A. A., P. Osland, A. A. Pankov, N. V. Romanenko, and J. Sirkka, 1998, Phys. Rev. D **58**, 075001.

Anoka, O. C., K. S. Babu, and I. Gogoladze, 2004, Nucl. Phys. B **687**, 3.

Antoniadis, I., 1990, Phys. Lett. B **246**, 377.

Appelquist, T., H.-C. Cheng, and B. A. Dobrescu, 2001, Phys. Rev. D **64**, 035002.

Appelquist, T., B. A. Dobrescu, and A. R. Hopper, 2003, Phys. Rev. D **68**, 035012.

Appelquist, T., and H.-U. Yee, 2003, Phys. Rev. D **67**, 055002.

Aranda, A., and C. D. Carone, 1998, Phys. Lett. B **443**, 352.

Arhrib, A., K. Cheung, C.-W. Chiang, and T.-C. Yuan, 2006, Phys. Rev. D **73**, 075015.

Arkani-Hamed, N., A. G. Cohen, and H. Georgi, 2001, Phys. Lett. B **513**, 232.

Arkani-Hamed, N., A. G. Cohen, E. Katz, and A. E. Nelson, 2002, J. High Energy Phys. **07**, 034.

Arkani-Hamed, N., and S. Dimopoulos, 2005, J. High Energy Phys. **06**, 073.

Arkani-Hamed, N., L. J. Hall, H. Murayama, D. R. Smith, and N. Weiner, 2001, Phys. Rev. D **64**, 115011.

Armillis, R., C. Coriano, and M. Guzzi, 2008, J. High Energy Phys. **05**, 015.

Babich, A. A., A. A. Pankov, and N. Paver, 1999, Phys. Lett. B **452**, 355.

Babu, K. S., X.-G. He, and E. Ma, 1987, Phys. Rev. D **36**, 878.

Babu, K. S., C. F. Kolda, and J. March-Russell, 1996, Phys. Rev. D **54**, 4635.

Babu, K. S., C. F. Kolda, and J. March-Russell, 1998, Phys. Rev. D **57**, 6788.

Baek, S., J. H. Jeon, and C. S. Kim, 2006, Phys. Lett. B **641**, 183.

Barbieri, R., A. Pomarol, R. Rattazzi, and A. Strumia, 2004, Nucl. Phys. B **703**, 127.

Barger, V., C.-W. Chiang, J. Jiang, and P. Langacker, 2004, Phys. Lett. B **596**, 229.

Barger, V., C.-W. Chiang, P. Langacker, and H.-S. Lee, 2004a, Phys. Lett. B **598**, 218.

Barger, V., C.-W. Chiang, P. Langacker, and H.-S. Lee, 2004b, Phys. Lett. B **580**, 186.

Barger, V., C. Kao, P. Langacker, and H.-S. Lee, 2004, Phys. Lett. B **600**, 104.

Barger, V., P. Langacker, and H.-S. Lee, 2003, Phys. Rev. D **67**, 075009.

Barger, V., P. Langacker, and H.-S. Lee, 2005, Phys. Lett. B **630**, 85.

Barger, V., P. Langacker, H.-S. Lee, and G. Shaughnessy, 2006, Phys. Rev. D **73**, 115010.

Barger, V., P. Langacker, and G. Shaughnessy, 2007a, Phys. Rev. D **75**, 055013.

Barger, V., P. Langacker, and G. Shaughnessy, 2007b, Phys. Lett. B **644**, 361.

Barger, V., P. Langacker, and G. Shaughnessy, 2007c, New J. Phys. **9**, 333.

Barger, V., *et al.*, 2007d, Phys. Rev. D **75**, 115002.

Barger, V. D., N. Deshpande, R. J. N. Phillips, and K. Whisnant, 1986, Phys. Rev. D **33**, 1912.

Barger, V. D., N. G. Deshpande, J. L. Rosner, and K. Whisnant, 1987, Phys. Rev. D **35**, 2893.

Barger, V. D., W. Y. Keung, and E. Ma, 1980, Phys. Rev. D **22**, 727.

Barger, V. D., E. Ma, and K. Whisnant, 1982a, Phys. Rev. Lett. **48**, 1589.

Barger, V. D., E. Ma, and K. Whisnant, 1982b, Phys. Rev. D **26**, 2378.

Barger, V. D., and R. J. N. Phillips, 1978, Phys. Rev. D **18**, 775.

Barger, V. D., and K. Whisnant, 1987, Phys. Rev. D **36**, 3429.

Barr, S. M., 1983, Phys. Lett. B **128**, 400.

Barr, S. M., B. Bednarz, and C. Benesh, 1986, Phys. Rev. D **34**, 235.

Barr, S. M., and A. Zee, 1980, Phys. Lett. B **92**, 297.

Batra, P., B. A. Dobrescu, and D. Spivak, 2006, J. Math. Phys. **47**, 082301.

Baumgart, M., T. Hartman, C. Kilic, and L.-T. Wang, 2007, J. High Energy Phys. **11**, 084.

Baur, U., D. Schildknecht, and K. H. G. Schwarzer, 1987, Phys. Rev. D **35**, 297.

Belanger, G., A. Pukhov, and G. Servant, 2008, J. Cosmol. Astropart. Phys. **0801**, 009.

Berenstein, D., and S. Pinansky, 2007, Phys. Rev. D **75**, 095009.
Binetruy, P., and P. Ramond, 1995, Phys. Lett. B **350**, 49.
Blumenhagen, R., M. Cvetic, P. Langacker, and G. Shiu, 2005, Annu. Rev. Nucl. Part. Sci. **55**, 71.
Blumenhagen, R., B. Kors, D. Lust, and S. Stieberger, 2007, Phys. Rep. **445**, 1.
Borzumati, F., G. R. Farrar, N. Polonsky, and S. D. Thomas, 1999, Nucl. Phys. B **555**, 53.
Borzumati, F., and Y. Nomura, 2001, Phys. Rev. D **64**, 053005.
Braun, V., Y.-H. He, B. A. Ovrut, and T. Pantev, 2005, J. High Energy Phys. **06**, 039.
Cacciapaglia, G., C. Csaki, G. Marandella, and A. Strumia, 2006, Phys. Rev. D **74**, 033011.
Carena, M. S., A. Daleo, B. A. Dobrescu, and T. M. P. Tait, 2004, Phys. Rev. D **70**, 093009.
Carena, M. S., A. Delgado, E. Ponton, T. M. P. Tait, and C. E. M. Wagner, 2003, Phys. Rev. D **68**, 035010.
Carena, M. S., M. Quiros, M. Seco, and C. E. M. Wagner, 2003, Nucl. Phys. B **650**, 24.
Carone, C. D., and H. Murayama, 1995, Phys. Rev. Lett. **74**, 3122.
Casalbuoni, R., S. De Curtis, D. Dominici, and R. Gatto, 1985, Phys. Lett. B **155**, 95.
Casalbuoni, R., S. De Curtis, D. Dominici, and R. Gatto, 1987, Nucl. Phys. B **282**, 235.
Casalbuoni, R., S. De Curtis, D. Dominici, and R. Gatto, 1999, Phys. Lett. B **462**, 48.
Chacko, Z., H.-S. Goh, and R. Harnik, 2006, Phys. Rev. Lett. **96**, 231802.
Chamseddine, A. H., and H. K. Dreiner, 1995, Nucl. Phys. B **447**, 195.
Chang, W.-F., J. N. Ng, and J. M. S. Wu, 2006, Phys. Rev. D **74**, 095005.
Chankowski, P. H., K. Kowalska, S. Lavignac, and S. Pokorski, 2005, Phys. Rev. D **71**, 055004.
Chankowski, P. H., S. Pokorski, and J. Wagner, 2006, Eur. Phys. J. C **47**, 187.
Chen, C.-H., and H. Hatanaka, 2006, Phys. Rev. D **73**, 075003.
Chen, M.-C., A. de Gouvea, and B. A. Dobrescu, 2007, Phys. Rev. D **75**, 055009.
Cheng, H.-C., B. A. Dobrescu, and K. T. Matchev, 1998, Phys. Lett. B **439**, 301.
Cheng, H.-C., B. A. Dobrescu, and K. T. Matchev, 1999, Nucl. Phys. B **543**, 47.
Cheng, H.-C., and I. Low, 2003, J. High Energy Phys. **09**, 051.
Cheng, H.-C., K. T. Matchev, and M. Schmaltz, 2002, Phys. Rev. D **66**, 056006.
Cheung, K., C.-W. Chiang, N. G. Deshpande, and J. Jiang, 2007, Phys. Lett. B **652**, 285.
Cheung, K.-m., 2001, Phys. Lett. B **517**, 167.
Cheung, K.-m., and G. L. Landsberg, 2002, Phys. Rev. D **65**, 076003.
Chiang, C.-W., N. G. Deshpande, and J. Jiang, 2006, J. High Energy Phys. **08**, 075.
Chivukula, R. S., H.-J. He, J. Howard, and E. H. Simmons, 2004, Phys. Rev. D **69**, 015009.
Chivukula, R. S., and E. H. Simmons, 2002, Phys. Rev. D **66**, 015006.
Cho, G.-C., K. Hagiwara, and Y. Umeda, 1998, Nucl. Phys. B **531**, 65.
Choi, S. Y., H. E. Haber, J. Kalinowski, and P. M. Zerwas, 2007, Nucl. Phys. B **778**, 85.
Chun, E. J., 2005, Phys. Rev. D **72**, 095010.
Cleaver, G., M. Cvetic, J. R. Espinosa, L. L. Everett, and P. Langacker, 1998, Phys. Rev. D **57**, 2701.
Cleaver, G., *et al.*, 1999, Phys. Rev. D **59**, 055005.
Coriano, C., A. E. Faraggi, and M. Guzzi, 2008, Eur. Phys. J. C **53**, 421.
Coriano, C., N. Irges, and E. Kiritsis, 2006, Nucl. Phys. B **746**, 77.
Costa, G., J. R. Ellis, G. L. Fogli, D. V. Nanopoulos, and F. Zwirner, 1988, Nucl. Phys. B **297**, 244.
Csaki, C., C. Grojean, L. Pilo, and J. Terning, 2004, Phys. Rev. Lett. **92**, 101802.
Cui, Y., S. P. Martin, D. E. Morrissey, and J. D. Wells, 2008, Phys. Rev. D **77**, 043528.
Cvetic, M., D. A. Demir, J. R. Espinosa, L. L. Everett, and P. Langacker, 1997, Phys. Rev. D **56**, 2861.
Cvetic, M., and S. Godfrey, 1995, e-print arXiv:hep-ph/9504216.
Cvetic, M., and P. Langacker, 1992a, Phys. Rev. D **46**, R14.
Cvetic, M., and P. Langacker, 1992b, Phys. Rev. D **46**, 4943.
Cvetic, M., and P. Langacker, 1996a, Phys. Rev. D **54**, 3570.
Cvetic, M., and P. Langacker, 1996b, Mod. Phys. Lett. A **11**, 1247.
Cvetic, M., and P. Langacker, 1997, e-print arXiv:hep-ph/9707451.
Cvetic, M., P. Langacker, and B. Kayser, 1992, Phys. Rev. Lett. **68**, 2871.
Cvetic, M., P. Langacker, and G. Shiu, 2002, Phys. Rev. D **66**, 066004.
Cvetic, M., G. Shiu, and A. M. Uranga, 2001, Nucl. Phys. B **615**, 3.
Davidson, A., 1979, Phys. Rev. D **20**, 776.
Davidson, S., S. Forte, P. Gambino, N. Rius, and A. Strumia, 2002, J. High Energy Phys. **02**, 037.
de Carlos, B., and J. R. Espinosa, 1997, Phys. Lett. B **407**, 12.
Degrassi, G., and A. Sirlin, 1989, Phys. Rev. D **40**, 3066.
de Groot, E. H., G. J. Gounaris, and D. Schildknecht, 1980, Z. Phys. C: Part. Fields **5**, 127.
del Aguila, F., 1994, Acta Phys. Pol. B **25**, 1317.
del Aguila, F., and J. A. Aguilar-Saavedra, 2007, J. High Energy Phys. **11**, 072.
del Aguila, F., G. A. Blair, M. Daniel, and G. G. Ross, 1987, Nucl. Phys. B **283**, 50.
del Aguila, F., G. D. Coughlan, and M. Quiros, 1988, Nucl. Phys. B **307**, 633.
del Aguila, F., and M. Cvetic, 1994, Phys. Rev. D **50**, 3158.
del Aguila, F., M. Cvetic, and P. Langacker, 1993, Phys. Rev. D **48**, 969.
del Aguila, F., M. Cvetic, and P. Langacker, 1995, Phys. Rev. D **52**, 37.
del Aguila, F., W. Hollik, J. M. Moreno, and M. Quiros, 1992, Nucl. Phys. B **372**, 3.
del Aguila, F., M. Masip, and M. Perez-Victoria, 1995, Nucl. Phys. B **456**, 531.
del Aguila, F., and A. Mendez, 1981, Nucl. Phys. B **189**, 212.
del Aguila, F., J. M. Moreno, and M. Quiros, 1989, Phys. Rev. D **40**, 2481.
del Aguila, F., M. Quiros, and F. Zwirner, 1987, Nucl. Phys. B **284**, 530.
Delgado, A., A. Pomarol, and M. Quiros, 2000, J. High Energy Phys. **01**, 030.
Demir, D. A., and L. L. Everett, 2004, Phys. Rev. D **69**, 015008.
Demir, D. A., L. L. Everett, and P. Langacker, 2008, Phys. Rev. Lett. **100**, 091804.

Demir, D. A., G. L. Kane, and T. T. Wang, 2005, Phys. Rev. D **72**, 015012.
Deshpande, N. G., and D. Iskandar, 1980, Nucl. Phys. B **167**, 223.
Deshpande, N. G., and J. Trampetic, 1988, Phys. Lett. B **206**, 665.
Dienes, K. R., C. F. Kolda, and J. March-Russell, 1997, Nucl. Phys. B **492**, 104.
Dittmar, M., 1997, Phys. Rev. D **55**, 161.
Dittmar, M., A.-S. Nicollerat, and A. Djouadi, 2004, Phys. Lett. B **583**, 111.
Dobrescu, B. A., 1997, Phys. Lett. B **403**, 285.
Dobrescu, B. A., 2005, Phys. Rev. Lett. **94**, 151802.
Dobrescu, B. A., and I. Mocioiu, 2006, J. High Energy Phys. **11**, 005.
Donini, A., F. Feruglio, J. Matias, and F. Zwirner, 1997, Nucl. Phys. B **507**, 51.
Duncan, M. J., and P. Langacker, 1986, Nucl. Phys. B **277**, 285.
Durkin, L. S., and P. Langacker, 1986, Phys. Lett. B **166**, 436.
Dvali, G. R., and A. Pomarol, 1998, Nucl. Phys. B **522**, 3.
Ellis, J. R., J. F. Gunion, H. E. Haber, L. Roszkowski, and F. Zwirner, 1989, Phys. Rev. D **39**, 844.
Erler, J., 2000, Nucl. Phys. B **586**, 73.
Erler, J., and P. Langacker, 1999, Phys. Lett. B **456**, 68.
Erler, J., and P. Langacker, 2000, Phys. Rev. Lett. **84**, 212.
Erler, J., P. Langacker, and T.-j. Li, 2002, Phys. Rev. D **66**, 015002.
Everett, L. L., P. Langacker, M. Plumacher, and J. Wang, 2000, Phys. Lett. B **477**, 233.
Faraggi, A. E., 1993, Phys. Rev. D **47**, 5021.
Faraggi, A. E., and D. V. Nanopoulos, 1991, Mod. Phys. Lett. A **6**, 61.
Fayet, P., 1977, Phys. Lett. B **69**, 489.
Fayet, P., 1980, Phys. Lett. B **96**, 83.
Fayet, P., 2007, Phys. Rev. D **75**, 115017.
Feldman, D., Z. Liu, and P. Nath, 2007, in *Theoretical High Energy Physics*, edited by A. Misra, AIP Conf. Proc. No. 939 (AIP, Melville, NY), p. 50.
Fiandrino, A., and P. Taxil, 1992, Phys. Lett. B **293**, 242.
Foot, R., and X.-G. He, 1991, Phys. Lett. B **267**, 509.
Foot, R., X. G. He, H. Lew, and R. R. Volkas, 1994, Phys. Rev. D **50**, 4571.
Freitas, A., 2004, Phys. Rev. D **70**, 015008.
Fuks, B., M. Klasen, F. Ledroit, Q. Li, and J. Morel, 2008, Nucl. Phys. B **797**, 322.
Georgi, H., and S. L. Glashow, 1974, Phys. Rev. Lett. **32**, 438.
Georgi, H., E. E. Jenkins, and E. H. Simmons, 1990, Nucl. Phys. B **331**, 541.
Georgi, H., and S. Weinberg, 1978, Phys. Rev. D **17**, 275.
Gherghetta, T., T. A. Kaeding, and G. L. Kane, 1998, Phys. Rev. D **57**, 3178.
Ghilencea, D. M., L. E. Ibanez, N. Irges, and F. Quevedo, 2002, J. High Energy Phys. **08**, 016.
Giedt, J., 2001, Ann. Phys. **289**, 251.
Giudice, G. F., and A. Masiero, 1988, Phys. Lett. B **206**, 480.
Godfrey, S., 2002, e-print arXiv:hep-ph/0201093.
Godfrey, S., P. Kalyniak, and A. Tomkins, 2005, e-print arXiv:hep-ph/0511335.
Gogoladze, I., and C. Macesanu, 2006, Phys. Rev. D **74**, 093012.
Gogoladze, I., and A. Perez-Lorenzana, 2002, Phys. Rev. D **65**, 095011.
Gonzalez-Garcia, M. C., and J. W. F. Valle, 1991, Phys. Lett. B **259**, 365.
Ham, S. W., and S. K. Oh, 2007, Phys. Rev. D **76**, 095018.
Ham, S. W., E. J. Yoo, and S. K. Oh, 2007, Phys. Rev. D **76**, 075011.
Hambye, T., E. Ma, M. Raidal, and U. Sarkar, 2001, Phys. Lett. B **512**, 373.
Han, T., P. Langacker, and B. McElrath, 2004, Phys. Rev. D **70**, 115006.
Han, T., H. E. Logan, B. McElrath, and L.-T. Wang, 2003, Phys. Rev. D **67**, 095004.
Han, T., H. E. Logan, and L.-T. Wang, 2006, J. High Energy Phys. **01**, 099.
Han, T., Z. Si, K. M. Zurek, and M. J. Strassler, 2007, J. High Energy Phys. **0807**, 008.
Han, T., G. Valencia, and Y. Wang, 2004, Phys. Rev. D **70**, 034002.
He, H.-J., *et al.*, 2008, Phys. Rev. D **78**, 031701.
He, X.-G., and G. Valencia, 2006, Phys. Rev. D **74**, 013011.
Hesselbach, S., F. Franke, and H. Fraas, 2002, Eur. Phys. J. C **23**, 149.
Hewett, J. L., F. J. Petriello, and T. G. Rizzo, 2002, J. High Energy Phys. **09**, 030.
Hewett, J. L., and T. G. Rizzo, 1989, Phys. Rep. **183**, 193.
Hewett, J. L., and T. G. Rizzo, 1993, Phys. Rev. D **47**, 4981.
Hill, C. T., 1995, Phys. Lett. B **345**, 483.
Hill, C. T., and R. J. Hill, 2007, Phys. Rev. D **76**, 115014.
Hill, C. T., and E. H. Simmons, 2003, Phys. Rep. **381**, 235.
Holdom, B., 1986, Phys. Lett. B **166**, 196.
Howl, R., and S. F. King, 2008, J. High Energy Phys. **01**, 030.
Hubisz, J., P. Meade, A. Noble, and M. Perelstein, 2006, J. High Energy Phys. **01**, 135.
Hur, T., H.-S. Lee, and S. Nasri, 2008, Phys. Rev. D **77**, 015008.
Ibanez, L. E., and G. G. Ross, 1994, Phys. Lett. B **332**, 100.
Jain, V., and R. Shrock, 1995, Phys. Lett. B **352**, 83.
Joshipura, A. S., R. D. Vaidya, and S. K. Vempati, 2000, Phys. Rev. D **62**, 093020.
Kang, J., and P. Langacker, 2005, Phys. Rev. D **71**, 035014.
Kang, J., P. Langacker, T.-j. Li, and T. Liu, 2005, Phys. Rev. Lett. **94**, 061801.
Kang, J., P. Langacker, and B. D. Nelson, 2008, Phys. Rev. D **77**, 035003.
Kang, J.-h., P. Langacker, and T.-j. Li, 2005, Phys. Rev. D **71**, 015012.
Kaplan, D. E., F. Lepeintre, A. Masiero, A. E. Nelson, and A. Riotto, 1999, Phys. Rev. D **60**, 055003.
Kawasaki, M., K. Kohri, and T. Moroi, 2005, Phys. Rev. D **71**, 083502.
Keith, E., and E. Ma, 1996, Phys. Rev. D **54**, 3587.
Keith, E., and E. Ma, 1997, Phys. Rev. D **56**, 7155.
Kim, J. E., and H. P. Nilles, 1984, Phys. Lett. B **138**, 150.
King, S. F., S. Moretti, and R. Nevzorov, 2006, Phys. Rev. D **73**, 035009.
Kiritsis, E., 2004, Fortschr. Phys. **52**, 200.
Kolda, C. F., and S. P. Martin, 1996, Phys. Rev. D **53**, 3871.
Kors, B., and P. Nath, 2004, Phys. Lett. B **586**, 366.
Kraan, A. C., J. B. Hansen, and P. Nevski, 2007, Eur. Phys. J. C **49**, 623.
Kumar, J., A. Rajaraman, and J. D. Wells, 2008, Phys. Rev. D **77**, 066011.
Kumar, J., and J. D. Wells, 2006, Phys. Rev. D **74**, 115017.
Langacker, P., 1981, Phys. Rep. **72**, 185.
Langacker, P., 1984, Phys. Rev. D **30**, 2008.
Langacker, P., 1998, Phys. Rev. D **58**, 093017.

Langacker, P., and D. London, 1988, Phys. Rev. D **38**, 886.
Langacker, P., and M.-x. Luo, 1992, Phys. Rev. D **45**, 278.
Langacker, P., M.-x. Luo, and A. K. Mann, 1992, Rev. Mod. Phys. **64**, 87.
Langacker, P., and M. Plumacher, 2000, Phys. Rev. D **62**, 013006.
Langacker, P., N. Polonsky, and J. Wang, 1999, Phys. Rev. D **60**, 115005.
Langacker, P., R. W. Robinett, and J. L. Rosner, 1984, Phys. Rev. D **30**, 1470.
Langacker, P., and S. Uma Sankar, 1989, Phys. Rev. D **40**, 1569.
Langacker, P., and J. Wang, 1998, Phys. Rev. D **58**, 115010.
Langacker, P. G., G. Paz, L.-T. Wang, and I. Yavin, 2008, Phys. Rev. Lett. **100**, 041802.
Lebedev, O., *et al.*, 2008, Phys. Rev. D **77**, 046013.
Lee, H.-S., C. Luhn, and K. T. Matchev, 2007, J. High Energy Phys. **0807**, 065.
Lee, H.-S., K. T. Matchev, and S. Nasri, 2007, Phys. Rev. D **76**, 041302.
Lee, H.-S., K. T. Matchev, and T. T. Wang, 2008, Phys. Rev. D **77**, 015016.
Leike, A., 1999, Phys. Rep. **317**, 143.
Leike, A., and S. Riemann, 1997, Z. Phys. C **75**, 341.
Leroux, K., and D. London, 2002, Phys. Lett. B **526**, 97.
Li, X., and R. E. Marshak, 1982, Phys. Rev. D **25**, 1886.
London, D., and J. L. Rosner, 1986, Phys. Rev. D **34**, 1530.
Lopez, J. L., and D. V. Nanopoulos, 1997, Phys. Rev. D **55**, 397.
Luhn, C., and M. Thormeier, 2008, Phys. Rev. D **77**, 056002.
Ma, E., 1987, Phys. Rev. D **36**, 274.
Ma, E., 1996, Phys. Lett. B **380**, 286.
Ma, E., 2002, Phys. Rev. Lett. **89**, 041801.
Maldacena, J. M., 1998, Adv. Theor. Math. Phys. **2**, 231.
Marciano, W. J., and A. Sirlin, 1987, Phys. Rev. D **35**, 1672.
Martin, S. P., 1997, e-print arXiv:hep-ph/9709356.
Masiero, A., 1980, Phys. Lett. B **93**, 295.
Masip, M., and A. Pomarol, 1999, Phys. Rev. D **60**, 096005.
Matsuoka, T., and D. Suematsu, 1986, Prog. Theor. Phys. **76**, 901.
Menon, A., D. E. Morrissey, and C. E. M. Wagner, 2004, Phys. Rev. D **70**, 035005.
Mohapatra, R. N., 1986, *Unification and Supersymmetry* (Springer, Berlin).
Mohapatra, R. N., and A. Riotto, 1997, Phys. Rev. D **55**, 4262.
Mohapatra, R. N., and G. Senjanovic, 1980, Phys. Rev. Lett. **44**, 912.
Mohapatra, R. N., and D. P. Sidhu, 1978, Phys. Rev. D **18**, 856.
Mohapatra, R. N., *et al.*, 2007, Phys. Rep. **70**, 1757.
Morrissey, D. E., and J. D. Wells, 2006, Phys. Rev. D **74**, 015008.
Morrissey, D. E., and J. D. Wells, 2007, J. High Energy Phys. **01**, 102.
Murakami, B., and J. D. Wells, 2003, Phys. Rev. D **68**, 035006.
Nakamura, S., and D. Suematsu, 2007, Phys. Rev. D **75**, 055004.
Nakayama, Y., 2008, J. High Energy Phys. **02**, 013.
Nelson, A. E., and J. Walsh, 2008, Phys. Rev. D **77**, 033001.
Olive, K. A., D. N. Schramm, and G. Steigman, 1981, Nucl. Phys. B **180**, 497.
Panagiotakopoulos, C., and K. Tamvakis, 1999, Phys. Lett. B **469**, 145.
Perelstein, M., 2007, Prog. Part. Nucl. Phys. **58**, 247.
Peskin, M. E., and T. Takeuchi, 1990, Phys. Rev. Lett. **65**, 964.
Pospelov, M., A. Ritz, and M. B. Voloshin, 2008, Phys. Lett. B **662**, 53.
Profumo, S., M. J. Ramsey-Musolf, and G. Shaughnessy, 2007, J. High Energy Phys. **08**, 010.
Randall, L., and R. Sundrum, 1999, Phys. Rev. Lett. **83**, 3370.
Richard, F., 2003, e-print arXiv:hep-ph/0303107.
Rizzo, T. G., 1980, Phys. Rev. D **21**, 1214.
Rizzo, T. G., 1987, Phys. Lett. B **192**, 125.
Rizzo, T. G., 1991, Phys. Rev. D **44**, 202.
Rizzo, T. G., 1993a, Phys. Rev. D **48**, 4236.
Rizzo, T. G., 1993b, Phys. Rev. D **47**, 956.
Rizzo, T. G., 2006, e-print arXiv:hep-ph/0610104.
Rizzo, T. G., 2007, J. High Energy Phys. **05**, 037.
Rizzo, T. G., and G. Senjanovic, 1981, Phys. Rev. D **24**, 704.
Robinett, R. W., and J. L. Rosner, 1982a, Phys. Rev. D **26**, 2396.
Robinett, R. W., and J. L. Rosner, 1982b, Phys. Rev. D **25**, 3036.
Rosner, J. L., 1996a, Phys. Rev. D **54**, 1078.
Rosner, J. L., 1996b, Phys. Lett. B **387**, 113.
Stueckelberg, E. C. G., 1938, Helv. Phys. Acta **11**, 225.
Suematsu, D., 1998, Phys. Rev. D **57**, 1738.
Suematsu, D., 1999, Phys. Rev. D **59**, 055017.
Suematsu, D., and Y. Yamagishi, 1995, Int. J. Mod. Phys. A **10**, 4521.
Umeda, Y., G.-C. Cho, and K. Hagiwara, 1998, Phys. Rev. D **58**, 115008.
Verlinde, H., L.-T. Wang, M. Wijnholt, and I. Yavin, 2008, J. High Energy Phys. **02**, 082.
Weiglein, G., *et al.*, (LHC/LC Study Group), 2006, Phys. Rep. **426**, 47.
Weinberg, S., 1996, *The Quantum Theory of Fields* (Cambridge University Press, Cambridge).
Witten, E., 1985, Nucl. Phys. B **258**, 75.
Witten, E., 2001, e-print arXiv:hep-ph/0201018.
Yao, W. M., *et al.*, 2006, J. Phys. G **33**, 1.
Zee, A., and J. E. Kim, 1980, Phys. Rev. D **21**, 1939.