Local Discrete Symmetries
From Superstring Derived Models

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Abstract

Discrete and global symmetries play an essential role in many extensions of the Standard Model, for example, to preserve the proton lifetime, to prevent flavor changing neutral currents, etc. An important question is how can such symmetries survive in a theory of quantum gravity, like superstring theory. In a specific string model I illustrate how local discrete symmetries may arise in string models and play an important role in preventing fast proton decay and flavor changing neutral currents. The local discrete symmetry arises due to the breaking of the non–Abelian gauge symmetries by Wilson lines in the superstring models and forbids, for example dimension five operators which mediate rapid proton decay, to all orders of nonrenormalizable terms. In the context of models of unification of the gauge and gravitational interactions, it is precisely this type of local discrete symmetries that must be found in order to insure that a given model is not in conflict with experimental observations.

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Discrete and global symmetries play a crucial role in many extensions of the Standard Model. Imposing such symmetries is in general necessary to insure agreement with various experimental observations. One example is the proton lifetime in the context of supersymmetric and superstring theories [1, 2, 3]. Supersymmetric theories give rise to dimension four and five operators which may result in rapid proton decay. Forbidding such operators requires that we impose some discrete or global symmetry on the spectrum of specific models. Another example is the flavor changing neutral currents in supersymmetric models which requires the flavor degeneracy of the soft breaking scalar masses. For example, in models of low–energy dynamical SUSY breaking, the supersymmetry breaking is mediated to the observable sector by a messenger sector which consists of down–like quarks and electroweak doublets. The SUSY breaking is mediated to the observable sector by the gauge interactions of the Standard Model, which are flavor blind. These messenger sector states would in general have flavor dependent interactions with the Standard Model quarks which will induce flavor changing neutral currents. It is therefore imperative that we impose a discrete symmetry which prevents the undesired interactions.

In the framework of point quantum field theories, it is of course simple to impose such discrete and global symmetries. However, it is well known that quantum gravity effects are, in general, expected to violate global and discrete symmetries [4]. The only exception to this expectation are local discrete symmetries [5]. Local discrete symmetries are discrete symmetries which arise from broken gauge symmetry. However, it is still difficult to envision how such symmetries will arise from a fundamental theory of gravity. The problem is best illustrated in the context of the realistic free fermionic superstring models [6, 7, 8]. In these models the cubic level and higher order terms in the superpotential are obtained by evaluating the correlators between the vertex operators [9, 10]

$$A_N \sim \langle V_1^f V_2^f V_3^b V_4^b \cdots V_N^b \rangle, \tag{1}$$

where V_i^f (V_i^b) are the fermionic (scalar) components of the vertex operators. The non–vanishing terms are obtained by applying the rules of Ref. [10].

2

The realistic free fermionic models contain an anomalous $U(1)$ symmetry. The anomalous $U(1)$ generates a Fayet–Iliopoulos term which breaks supersymmetry and destabilizes the vacuum [11]. Supersymmetry is restored and the vacuum is stabilized by assigning VEVs to a set of Standard Model singlets in the massless string spectrum, which break the anomalous $U(1)$ symmetry. These Standard Model singlets in general also carry charges under the non–anomalous $U(1)$ symmetries which exist in the superstring models. Therefore, requiring that all the D–terms and F–terms vanish imposes a set of non–trivial constraints on the allowed VEVs. In this process some of the fields in the higher order nonrenormalizable terms in Eq. (1) accrue a VEV. Some of the nonrenormalizable terms then become effective renormalizable operators. These VEVs, in general, will also break most or all of the additional local $U(1)$ symmetries and the global and discrete symmetries. So, although, some terms may be forbidden up to some order in the superpotential, it is difficult to envision, how, in general, a term which is not protected by an unbroken local symmetry will not be generated at some order [12]. However, in several phenomenological cases, the experimental constraints are so severe that we must insure that the dangerous terms are forbidden up to a very high order. For example, this is the case with regard to the problems of proton stability and FCNC in supersymmetric theories. For instance, if we assume that each VEV produces a suppression factor of order 1/10 then to insure that dimension four baryon and lepton violating operators are not induced up to order $N = 14 - 15$. In practice, one finds that in general, the dangerous operators are induced at various orders [12]. If the suppression of some of the singlets VEVs is larger, or perhaps even of order one, then one has to go to even higher orders to insure agreement with the experimental data.

In these paper I discuss how such phenomenologically disastrous operators may be avoided in superstring models to all orders of nonrenormalizable terms. The symmetry which forbids the undesired operators arises as follows. The free fermionic models correspond to orbifold models of toroidally compactified models [15]. In These models we start with a large symmetry group, like $SO(44)$ or $SO(12) \times E_8 \times E_8$ or $SO(12) \times SO(16) \times SO(16)$, and with $N = 4$ supersymmetry. The number of super-

symmetries is reduced to one and the gauge group is broken to one of its subgroup by the orbifolding. In the realistic free fermionic models, the $SO(12) \times SO(16) \times SO(16)$ is typically broken to $SO(4)^3 \times SO(10) \times U(1)^3 \times SO(16)$. Alternative three generation free fermionic models starting with an $SO(44)$ gauge group were discussed in ref. [13]. The $SO(10)$ symmetry is then broken further to one of its subgroups by additional boundary condition basis vectors. These additional boundary condition basis vectors correspond to Wilson lines in the orbifold formulation. The breaking of the non-Abelian gauge symmetries by Wilson lines gives rise to massless states that do not fall into representations of the original unbroken $SO(10)$ symmetry. This is an intrinsic stringy phenomena. I refer to the states from these sectors as Wilsonian matter states. The basis vectors which break the $SO(10)$ gauge symmetry generate sectors in the partition function which break the $SO(10)$ symmetry. The massless states from these sectors carry fractional charges under the $U(1)$ symmetries which are embedded in $SO(10)$ and which are orthogonal to the generators of $SU(3) \times SU(2)$. Thus, they can carry fractional charge under $U(1)_Y$, the weak hypercharge, or they can carry fractional charge under $U(1)_{Z'}$ which is embedded in $SO(10)$ and is orthogonal to the weak hypercharge. The charge of the Wilsonian states under these $U(1)$ symmetries, do not have the standard $SO(10)$ quantization. Because of the appearance of this type of matter from the Wilsonian sectors, we can now get conservations laws which forbid the interactions of the Wilsonian states with the Standard Model states. For example, if the $U(1)_{Z'}$ symmetry is broken only by a VEV of the right-handed neutrino then there will be a residual discrete symmetry which forbids the coupling of the Wilsonian matter states to the Standard Model states.

The emergence of local discrete symmetries in superstring models is nicely illustrated in the model of ref. [8]. This model is constructed in the free fermionic formulation [14] and belongs to a subset of free fermionic models. These models utilize a set of boundary condition basis vectors which correspond to $Z_2 \times Z_2$ orbifold compactification with standard embedding [15]. To show how the local discrete symmetry appears in the string model, I first discuss the general structure of this class of free fermionic models. More details on the construction of the realistic free fermionic

4

models are given in ref. [7].

In the free fermionic formulation all the degrees of freedom which are needed to cancel the conformal anomaly are represented in terms of free fermions propagating on the string world–sheet. Under parallel transport around one of the noncontractible loops of the world–sheet torus, the fermionic states pick up a phase. These phases for all the 64 world–sheet fermions are collected in the diagonal boundary condition basis vectors. and span a finite additive group. Modular transformation in general mix between the different spin structures. Requiring invariance under the modular transformations restricts the possible choices of boundary condition basis vectors and the one–loop phases. The physical spectrum is obtained by applying the generalized GSO projections. The quantum numbers with respect to the Cartan generators of the four dimensional gauge group are given by

$$Q(f) = \frac{1}{2}\alpha(f) \; + \; F(f) \tag{2}$$

where f is a complex world–sheet fermion which produces a space–time $U(1)$ current of the four dimensional gauge group, $\alpha(f)$ and $F(f)$ are the boundary condition and fermion number of the fermion f in the sector α. Each state in the physical spectrum and its charges under the four dimensional gauge are represented in terms of a vertex operator. The cubic level and higher order terms in the superpotential are obtained by evaluating the correlators between the vertex operators. Following this procedure we can construct the string physical spectrum and study its phenomenology.

The basis which generate the realistic free fermionic models, typically consists of eight or nine boundary condition basis vectors. These are typically denoted by $\{1, S, b_1, b_2, b_3, \alpha, \beta, \gamma\}$. The boundary conditions correspond to orbifold twisting and Wilson lines in the corresponding bosonic construction. However the correspondence is usually not apparent. The construction of the free fermionic standard–like models is divided to two parts. The first part consist of the boundary condition basis vector of the NAHE set, $\{1, S, b_1, b_2, b_3\}$. This set of boundary condition basis vectors plus the basis vector 2γ, correspond to $Z_2 \times Z_2$ orbifold compactification with standard embedding. The set $\{1, S, \xi = 1 + b_1 + b_2 + b_3, 2\gamma\}$ produces a toroidally compactified model with $N = 4$ space–time supersymmetry and $SO(12) \times SO(16) \times SO(16)$ gauge

group. The action of the basis vectors b_1 and b_2 corresponds to the $Z_2 \times Z_2$ twisting and reduce the number of supersymmetries to $N = 1$ and the gauge group is broken to $SO(10) \times U(1)^3 \times SO(16) \times SO(4)^3 \times$. The NAHE set plus the vector 2γ is common to a large number of realistic free fermionic models and to all the models which are discussed in this paper. At this level each one of the basis vectors b_1, b_2 and b_3 gives rise to eight generations in the chiral 16 representation of $SO(10)$.

The next stage in the construction of the realistic free fermionic models is the construction of the basis vectors $\{\alpha, \beta, \gamma\}$. This set of boundary condition basis vectors reduces the number of generations to three generations one from each of the sectors b_1, b_2 and b_3. At the same time the $SO(10)$ gauge group is broken to one of its subgroups, $SO(6) \times SO(4)$, $SU(5) \times U(1)$ or $SU(3) \times SU(2) \times U(1)^2$. The hidden $SO(16)$ is also broken to one of its subgroups and the horizontal $SO(6)^3$ symmetries are broken to $U(1)^n$, where n can vary between three and nine. In the free fermionic standard–like models the $SO(10)$ symmetry is broken to $SU(3) \times SU(2) \times U(1)_C \times U(1)_L{}^*$. The weak hypercharge is given by

$$U(1)_Y = 1/3 U(1)_C + 1/2 U(1)_L$$

and the orthogonal $U(1)_{Z'}$ combination is given by

$$U(1)_{Z'} = U(1)_C - U(1)_L.$$

The three twisted sectors b_1, b_2 and b_3 produce three generations in the sixteen representation of $SO(10)$ decomposed under the final $SO(10)$ subgroup. These states carry half integral charges under the $U(1)_{Z'}$ gauge symmetry,

$$e_L^c \equiv [(1, \frac{3}{2}); (1, 1)]_{(1,1/2,1)}; \tag{3}$$

$$u_L^c \equiv [(\bar{3}, -\frac{1}{2}); (1, -1)]_{(-2/3,1/2,-2/3)}; \tag{4}$$

$$Q \equiv [(3, \frac{1}{2}); (2, 0)]_{(1/6,1/2,(2/3,-1/3))} \tag{5}$$

$$N_L^c \equiv [(1, \frac{3}{2}); (1, -1)]_{(0,5/2,0)}; \tag{6}$$

$^*U(1)_C = 3/2 U(1)_{B-L}$; $U(1)_L = 2 U(1)_{T_{3_R}}$

$$d_L^c \equiv [(\bar{3}, -\frac{1}{2}); (1,1)]_{(1/3,-3/2,1/3)}; \tag{7}$$

$$L \equiv [(1, -\frac{3}{2}); (2,0)]_{(-1/2,-3/2,(0,1))}, \tag{8}$$

where I have used the notation

$$[(SU(3)_C; U(1)_C); (SU(2); U(1)_L)]_{\{Q_Y, Q_{Z'}, Q_{e.m.}\}} \tag{9}$$

Similarly, the states which are identified with the light Higgs representations are obtained from $SO(10)$ representations which are broken by the GSO projections of the additional basis vectors α, β, γ.

The basis vectors α, β and γ correspond to Wilson lines in the bosonic formulation. These additional basis vectors give rise to additional massless spectrum. The massless states which arise due to the Wilson line breaking cannot fit into representations of the original unbroken $SO(10)$ symmetry. I will refer to these generically as exotic Wilsonian matter states. They carry non–standard $SO(10)$ charges under the $U(1)$ symmetries which are embedded in $SO(10)$. These two $U(1)$ are the weak–hypercharge, $U(1)_Y$, and an orthogonal combination, $U(1)_{Z'}$. Thus, the exotic Wilsonian states carry fractional charges under, $U(1)_Y$ or under $U(1)_{Z'}$.

Each Wilsonian sector in the additive group breaks the $SO(10)$ symmetry to one of its subgroups, $SO(6) \times SO(4)$, $SU(5) \times U(1)$ or $SU(3) \times SU(2) \times U(1)^2$. Thus, the physical states from each of these sectors are classified according to the pattern of $SO(10)$ symmetry breaking. Below I list all the exotic Wilsonian states which appear in the realistic free fermionic models and classify the states according to the pattern of symmetry breaking.

The $SO(6) \times SO(4)$ type sectors are sectors with boundary conditions $\{1,1,1,0,0\}$ for the complex fermions $\bar{\psi}^{1,\cdots,5}$. These type of sectors give rise to states with the charges

$$[(3, \frac{1}{2}); (1,0)]_{(1/6,1/2,1/6)} \quad ; \tag{10}$$

$$[(\bar{3}, -\frac{1}{2}); (1,0)]_{(-1/6,-1/2,-1/6)} ; \tag{11}$$

$$[(1,0); (2,0)]_{(0,0,\pm1/2)} ; \tag{12}$$

$$[(1,0);(1,\pm1)]_{(\pm1/2,\mp1/2,\pm1/2)} \tag{13}$$

$$[(1,\pm3/2);(1,0)]_{(\pm1/2,\pm1/2,\pm1/2)} \tag{14}$$

The $SU(5) \times U(1)$ type sectors are sectors with boundary conditions $\{1/2,1/2,1/2,1/2,1/2\}$ for the complex fermions $\bar{\psi}^{1,\cdots,5}$. These type of sectors give rise to states with the charges

$$[(1,\pm3/4);(1,\pm1/2)]_{(\pm1/2,\pm1/4,\pm1/2)} \tag{15}$$

Finally, the $SU(3) \times SU(2) \times U(1)^2$ type sectors are sectors with boundary conditions $\{1/2,1/2,1/2,-1/2,-1/2\}$ for the complex fermions $\bar{\psi}^{1,\cdots,5}$. These type of sectors give rise to states which carry the usual charges under the Standard Model gauge group but carry fractional charges under the $U(1)_{Z'}$ symmetry.

$$\begin{aligned}
&[(3,\tfrac{1}{4});(1,\tfrac{1}{2})]_{(-1/3,-1/4,-1/3)}; \\
&[(\bar{3},-\tfrac{1}{4});(1,\tfrac{1}{2})]_{(1/3,1/4,1/3)} ; \\
&[(1,\pm\tfrac{3}{4});(2,\pm\tfrac{1}{2})]_{(\pm1/2,\pm1/4,(1,0);(0,-1))} ; \\
&[(1,\pm\tfrac{3}{4});(1,\mp\tfrac{1}{2})]_{(0,\pm5/4,0)} \tag{16}
\end{aligned}$$

The $SO(6) \times SO(4)$ and $SU(5) \times U(1)$ type Wilsonian matter states carry fractional electric charge $\pm1/2$ and therefore must be either, confined, diluted or have a mass of the order of the Planck scale. Because the Wilsonian matter states appear in the realistic free fermionic models in vector–like representations, in general, they can get mass at a scale which is much higher than the electroweak scale. In specific string models, detailed scenarios were proposed in which these states are confined or become supermassive. The $SU(3) \times SU(2) \times U(1)^2$ type Wilsonian matter states transform as regular quarks and leptons under the Standard Model gauge group or are Standard Model singlets. These type of states may have important cosmological and phenomenological implications.

To illustrate how the local discrete symmetries arise in the the superstring models I focus on the Wilsonian color triplets in Eq. (16). These color triplets transform

under the Standard Model gauge group as right–handed down–type quarks, with weak hypercharge $\pm 1/3$. Thus, they can fit into the five representation of $SU(5)$. They may form interaction terms with the Standard Model states which are invariant under the Standard Model gauge group. However, they carry fractional charge under the $U(1)_{Z'}$ which is embedded in $SO(10)$. While the Standard Model states are obtained from the sectors b_1, b_2 and b_3 and have charges $n/2$ under the $U(1)_{Z'}$ symmetry, the Wilsonian color triplets have charges $\pm 1/4$ under the $U(1)_{Z'}$ symmetry. In Eq. (17) all the possible interaction terms of the Wilsonian triplets with the Standard Model states are written

$$LQ\bar{D}, \; u_L^c e_L^c D, \; QQD, \; u_L^c d_L^c \bar{D}, \; d_L^c N_L^c D,$$
$$QDh$$
$$\bar{D}\bar{D}u_L^c \tag{17}$$

The form of the interaction terms is $f_i f_j D \phi^n$ or $f_i D D \phi^n$ where f_i and f_j are the Standard Model states from the sectors b_1, b_2 and b_3 and D represents the Wilsonian triplet. The product of fields, ϕ^n, is a product of Standard Model singlets which insures invariance of the interaction terms under all the $U(1)$ symmetries and the string selection rules. If all the fields ϕ in the string ϕ^n get VEVs then the coefficients of the operators in Eq. (17) will be of the order $(\phi/M)^n$, where $M \sim 10^{18}$ GeV is a scale which is related to the string scale and I am assuming that the numerical coefficients of the correlators of the interactions terms are of order one. Because of the fractional charge of the Wilsonian color triples under the $U(1)_{Z'}$ all the interactions terms in Eq. (17) are not invariant under $U(1)_{Z'}$. The total $U(1)_{Z'}$ charge of each of these interaction terms is a multiple of $\pm(2n+1)/4$. Thus, for these terms to be allowed the string ϕ^n must break $U(1)_{Z'}$ and must must a total $U(1)_{Z'}$ charge in multiple of $\pm(2n+1)/4$. Thus, the string of Standard Model singlets must contain a field which carries fractional $U(1)_{Z'}$ charge $\pm(2n+1)/4$. In the model of ref. [8] the only Standard Model singlets with fractional $U(1)_{Z'}$ charge transform as triplets of the hidden $SU(3)_H$ gauge group. Therefore, if we make the single assumption that the hidden $SU(3)_H$ gauge group remains unbroken then all the interaction terms

9

between the Wilsonian triplet and the Standard Model states are suppressed to all orders of nonrenormalizable terms. In this case the $U(1)_{Z'}$ symmetry may be broken by the VEV of the right-handed sneutrino, which carry charge $Q_{Z'} = \pm 1/2$. Thus, in this case a residual Z_4 local discrete symmetry remains unbroken and suppresses the couplings of the Wilsonian triplets to the Standard Model states. However, since the states which transform under the hidden $SU(3)$ gauge group always appear in vector–like representations, invariance under the hidden $SU(3)$ guarantees that the discrete Z_4 symmetry remains unbroken also if the $U(1)_{Z'}$ gauge symmetry is broken by the VEVs of the hidden $SU(3)$ triplet representations. Thus, the local discrete Z_4 symmetry remains unbroken and forbids the couplings in Eq. (17) to all orders of nonrenormalizable terms.

The appearance of a good local discrete symmetry in this manner is an intriguing miracle. The phenomenological implications are striking. The string scale gauge coupling unification requires the existence of the Wilsonian color triplets at an intermediate energy scale [17]. However, the intermediate color triplets may, a priori, mediate rapid proton decay through dimension five operator. The existence of the local discrete symmetry forbids the dangerous dimension five operators. The existence of the local discrete symmetry indicates that the Wilsonian color triplets have interesting cosmological implications [18], and may result in testable experimental predictions of the superstring models. Finally, if we consider the color triplets as the messenger sector in dynamical SUSY breaking scenarios [19], then the local discrete symmetry guarantees that the interaction of the messenger sector with the Standard Model states occurs only through the gauge interactions. In this case indeed the problem with flavor changing neutral currents in supersymmetric models is resolved. In the context of models of unification of the gauge and gravitational interactions, it is precisely this type of local discrete symmetries that must be found in order to insure that a given model is not in conflict with experimental observations.

In this paper I have shown how local discrete symmetries may arise from superstring derived models. The proposed local discrete symmetries arise due to the breaking of the non–Abelian gauge symmetries by Wilson lines in the superstring

models. The breaking by Wilson lines give rise to massless states that cannot fit into representations of the original unbroken non–Abelian gauge symmetry while the Standard Model spectrum and phenomenology are obtained from representations of the original unbroken non–Abelian gauge symmetries. The unique stringy breaking of the non–Abelian gauge symmetries by Wilson lines may therefore result in local discrete symmetries which forbid the interactions of the Wilsonian matter states to the Standard Model states. The local discrete symmetries are good symmetries also when quantum gravity effects are taken into account and survive to all orders of nonrenormalizable terms. From the low energy point of view such local discrete symmetries are essential, for example, to prevent flavor changing neutral currents in gauge mediated dynamical SUSY breaking scenarios, to prevent rapid proton decay from dimension five operators, etc. The proposed local discrete symmetries were illustrated in a specific free fermionic model. However, the use of Wilson line breaking is common to a large class of superstring models. Therefore, similar symmetries may arise in other superstring standard–like models [20]. It will also be of interest to examine whether string models which do not use Wilson line breaking [21] give rise to similar symmetries. In the context of models of unification of the gauge and gravitational interactions, it is precisely this type of local discrete symmetries that must be found in order to insure that a given model is not in conflict with experimental observations.

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F	SEC	$SU(3)_C \times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	$SU(5)_H \times SU(3)_H$	Q_7	Q_8
D_1	b_2+b_3+	$(3,1)$	$\frac{1}{4}$	$\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	$(1,1)$	$-\frac{1}{4}$	$-\frac{15}{4}$
$\overline{D}_1$	$\beta+\gamma+\xi$	$(\overline{3},1)$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	0	$(1,1)$	$\frac{1}{4}$	$\frac{15}{4}$
D_2	b_1+b_3+	$(3,1)$	$\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	$(1,1)$	$\frac{1}{4}$	$-\frac{15}{4}$
$\overline{D}_2$	$\beta+\gamma+\xi$	$(\overline{3},1)$	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	0	0	0	$(1,1)$	$\frac{1}{4}$	$\frac{15}{4}$
H_1	b_2+b_3+	$(1,1)$	$-\frac{3}{4}$	$\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	$(1,3)$	$\frac{3}{4}$	$\frac{5}{4}$
$\overline{H}_1$	$\beta+\gamma+\xi$	$(1,1)$	$\frac{3}{4}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	0	$(1,\overline{3})$	$-\frac{3}{4}$	$\frac{5}{4}$
H_2	b_1+b_3+	$(1,1)$	$\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{3}{4}$	0	0	0	$(1,3)$	$\frac{3}{4}$	$-\frac{5}{4}$
$\overline{H}_2$	$\beta+\gamma+\xi$	$(1,1)$	$-\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{3}{4}$	0	0	0	$(1,\overline{3})$	$\frac{3}{4}$	$\frac{5}{4}$

Table 1: Massless Wilsonian states with fractional $U(1)_{Z'}$ charge in the model of ref. [8]. The first two pairs are the Wilsonian down–like color triplets. The last two pairs are the hidden sector triplets with vanishing weak hypercharge and fractional $U(1)_{Z'}$ charge.