Flavor universal dynamical electroweak symmetry breaking

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The top condensate seesaw mechanism of Dobrescu and Hill allows electroweak symmetry to be broken while deferring the problem of flavor to an electroweak singlet, massive sector. We provide an extended version of the singlet sector that naturally accommodates realistic masses for all the standard model fermions, which play an equal role in breaking electroweak symmetry. The models result in a relatively light composite Higgs sector with masses typically in the range of (400–700) GeV. In more complete models the dynamics will presumably be driven by a broken gauged family or flavor symmetry group. As an example of the higher scale dynamics a fully dynamical model of the quark sector with a GIM mechanism is presented, based on an earlier top condensation model of King using broken family gauge symmetry interactions (that model was itself based on a technicolor model of Georgi). The crucial extra ingredient is a reinterpretation of the condensates that form when several gauge groups become strong close to the same scale. A related technicolor model of Randall which naturally includes the leptons too may also be adapted to this scenario. We discuss the low energy constraints on the massive gauge bosons and scalars of these models as well as their phenomenology at the TeV scale. [S0556-2821(99)07609-2]

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I. INTRODUCTION

The electroweak symmetry (EWS) of the standard model, which is a chiral symmetry of the fermions, is spontaneously broken by some as yet unknown mechanism. The introduction of an elementary Higgs scalar is technically unnatural. A natural explanation for the breaking of chiral symmetry is a higher scale repeat of the dynamical breaking induced by QCD. This idea spawned technicolor models [1] of EWS breaking but these models ran into trouble since they typically introduce many extra electroweak doublet fermions whose presence is naively in conflict with precision experimental data from the CERN e^+e^- collider LEP and SLAC Linear Collider (SLC) [2]. With the discovery of the large top quark mass it was proposed that a top quark condensate [3], generated by some chiral symmetry breaking but not confining interaction (most likely a gauge group broken close to its critical scale for chiral symmetry breaking), might be responsible for EWS breaking. However, the value of the top quark mass is too small to generate the EWS breaking scale v. Introducing extra fermions in this fashion again leads to a conflict with electroweak precision measurements. A further problem, also found in technicolor, is that attempts to generate the top-bottom quark mass splitting typically give rise to large custodial isospin violating effects in the massive sector [4]. Such effects are harshly constrained by the precision measurements of the ρ parameter. A possible alternative is the combination of top condensation as the source of a large dynamical top quark mass and technicolor as giving most of the EWS breaking. In this scenario, known as top-color assisted technicolor [5], the extended technicolor interactions give the top quark only a small fraction of its mass. On the other hand, it is still necessary to have a large number of techni-doublets in order to obtain the correct pattern of light quark masses. A variety of models of this type as well as their phenomenological consequences have been studied in the literature [6].

Recently Dobrescu and Hill [7] have proposed a model in which a dynamically generated mass for the top quark does generate the full electroweak symmetry breaking scale but the correct top quark mass is obtained as a result of a seesaw mechanism with a heavy fermion sector. Perhaps of more significance, the model allows the origin of custodial symmetry breaking to be deferred to an EWS singlet sector where it cannot contaminate the ρ parameter. In their model the top quark is treated in a unique fashion, its condensation being driven by a broken color interaction unique to third generation quarks. This essential non-universality disrupts the $SU(3)$ family symmetry of the standard model (SM) in the absence of fermion masses. As a consequence, it is not obvious how to feed the top condensate down to provide masses for the leptons and lighter first and second generation quarks.

In this paper we propose a mechanism analogous to the top condensate seesaw mechanism but which can naturally accommodate all the fermion masses and generation structure. An interesting aspect of the resulting models is that all the SM electroweak doublets play an equal role in breaking EWS. The result of this universal involvement in EWS

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breaking is that the models give rise to a relatively light[1] composite Higgs sector with a typical mass scale of order (400–700) GeV. The SM fermion mass splittings are the result of mass terms in an EWS singlet sector, which at this stage of model building is not explained dynamically but simply put in by hand. In Sec. II we present the simplest version of the model in which the dynamical symmetry breaking is driven by strong four fermion interactions. Depending on the flavor structure of the four fermion interactions driving EWS breaking, different patterns of pseudo Nambu-Goldstone bosons (PNGBs) result. We discuss in detail one possible pattern in Sec. III which is analogous to the spectrum of a one family technicolor model, and another one in the model of Sec. IV where there are no unabsorbed PNGBs at the weak scale.

The simple model of Sec. II is in fact inspired by a more complete model of the dynamics of top condensation of King [8], which in turn was derived from a technicolor model by Georgi [9]. In Ref. [8] the four fermion interactions are the result of a broken gauged family symmetry. The biggest success of this class of models is that they have a Glashow-Iliopoolos-Maiani (GIM) mechanism above the weak scale which protects them from flavor changing neutral current (FCNC) even with family gauge bosons with masses of the order of 1 TeV. This is achieved by gauging the full chiral family symmetry, the symmetry responsible for the SM GIM mechanism. The model does have a number of possible flaws, including a non-trivial assumption about vacuum alignment and potentially light PNGBs coming from the singlet sector. Nevertheless, it is useful to elucidate the idea of a universal seesaw mechanism. The model is also suggestive of the gauge structure that is likely to underlie the universal seesaw model suggesting experimental searches. In Sec. III we revive the model by reinterpreting the dynamics when several gauge groups become strongly interacting at the same scale. The model then produces the flavor universal seesaw masses. The full model has additional interesting dynamics as a result of the massive, strongly interacting flavorons of the family gauge symmetry group and colorons analogous to those of the model of Ref. [10]. Their phenomenology, together with that of the scalar and pseudoscalar sector originating from the breaking of the fermionic chiral symmetries, is discussed. The inclusion of leptons in that model is non-trivial. A related technicolor model of Randall [11], which more naturally includes the leptons, may be adapted to provide a low energy flavor universal EWS breaking model where the strong four fermion interactions result from a larger gauged flavor symmetry. We discuss this model, its flaws and its phenomenology in Sec. IV.

Finally, we must address the matter of fine-tuning in top condensate models. In these models we will assume that the interactions responsible for EWS breaking are broken gauge interactions with gauge bosons with masses in the range $\Lambda = (1–10)$ TeV. The fact that their interactions give rise to the weak scale $v = 246$ GeV naively suggests fine-tuning of the order of (1–10)%. In fact in the underlying models we present, where the dynamics is more complete, a stronger degree of fine-tuning is probably required both because gauge couplings do not run linearly with momentum scale and because several gauge groups become strongly coupled at essentially the same scale. Although the reader may still find these tunings uncomfortable, they are clearly much less severe than those of the SM. We do not make any judgements and simply wish to explore this paradigm of model building. Thankfully in these matters we will eventually be instructed by experiment. One of the successes of the flavor universal models is that they allow direct condensation of all the electroweak doublets and hence masses for all the SM fermions without increasing the fine-tuning. In the original top condensation models, to generate the electron mass by a direct condensate would have involved fine-tuning of the order of m_e/Λ; the suppression of the mass in the present models is the result of the smallness of mass terms in a singlet sector. The structure also does not require very large singlet masses to seesaw the electron mass small which again would have been a source of fine-tuning since it would have driven the upper cutoff higher.

[1] We refer to Higgs boson masses in the few hundred GeV range as light, compared to the unitarity bound of approximately $\simeq 1.2$ TeV. A large class of models of a strongly coupled EWS breaking sector saturates this bound.

II. FLAVOR UNIVERSAL SEESAW MECHANISM

Our flavor universal seesaw model can be thought of as an extension of the top condensate seesaw model of Dobrescu and Hill [7], which we briefly review next. That model, in addition to the top quark doublet, $Q_L = (t_L, b_L)$ and the t_R contains the electroweak singlets χ_L, χ_R which have the same QCD and $U(1)_Y$ charges as the t_R. The EWS breaking is driven by the four fermion interaction

$$G\overline{Q}_L\chi_R\overline{\chi}_R Q_L, \tag{1}$$

which, if the coupling G is above critical, generates an EWS breaking mass between Q_L and χ_R. This can be seen in the large N approximation to the gap equation

$$\frac{1}{G} = \frac{N}{4\pi^2}\left[\Lambda^2 - m^2 \ln\left(\frac{\Lambda^2}{m^2}\right)\right], \tag{2}$$

where Λ is the scale above which the dynamics generating Eq. (1) resides, and m is the dynamically generated mass. In addition all EWS invariant masses are allowed, that is $\overline{t}_R\chi_L$ and $\overline{\chi}_R\chi_L$. The resulting mass matrix below the EWS breaking scale is then

$$(\overline{t_L}, \overline{\chi_L})\begin{pmatrix} 0 & m_{t_L\chi} \\ m_{t_R\chi} & m_{\chi\chi} \end{pmatrix}\begin{pmatrix} t_R \\ \chi_R \end{pmatrix}. \tag{3}$$

The value of $m_{t_L\chi}$ needed to provide all the EWS breaking vacuum expectation values (VEVs) may be estimated using the Pagels-Stokar formula [12]

$$v^2 \simeq \frac{N_c}{4\pi^2} m_{t_L\chi}^2 \ln\left(\frac{\Lambda}{m_{\chi\chi}}\right). \tag{4}$$

Thus, for instance for $\Lambda/m_{\chi\chi} \simeq O(10)$, $m_{t_L\chi} \approx 600$ GeV. By appropriate choices of the remaining two masses (e.g. $m_{t_R\chi} \simeq 1$ TeV and $m_{\chi\chi} \simeq 3$ TeV) the lightest mass eigenstate of the mass matrix comes out as 175 GeV, whereas the most massive is 3 TeV. In fact by suitable choices of the two singlet masses the value of the lightest eigenstate mass and the electroweak breaking mass may be maintained even in the limit of decoupling the χ fermion by taking $m_{\chi\chi} \rightarrow \infty$. Note that the electroweak breaking mass in that limit is only for the left handed top quark doublet and so does not violate custodial isospin.

There is a crucial difference between the top seesaw model and previous top quark condensate models. That is that the EWS breaking VEV of the left handed top is with an EWS singlet not the right handed top quark. The right handed top quark also has a mass with a singlet sector fermion. The SM top quark mass results from the mass mixing in the singlet sector indirectly connecting the left and right handed top quarks. The size of the top quark mass is no longer a direct consequence of its role in EWS breaking but simply of some singlet mass structure. The top quark is therefore in this sort of model no longer the essential fermion to be involved in EWS breaking. Any SM fermion could play the same role. In the seesaw mechanism of Eq. (3) the light mass eigenstate is essentially given by $m_{t_L\chi} m_{t_R\chi}/m_{\chi\chi}$. Using the same seesaw but generating e.g. the electron mass would require a very large value of $m_{\chi\chi}$. The upper cutoff on the theory would have to be raised and with it the degree of fine-tuning required to generate the weak scale. To avoid this we will enlarge the singlet sector. Consider for example the generation of the top quark mass. We introduce the additional electroweak singlet fields χ_L, χ_R but also ψ_L, ψ_R. We assume that these extra singlets are all colored under the usual SU(3) QCD group and have the same hypercharge as the t_R. The model then induces four fermion interactions of the form (suppressing color indices)

$$\bar{Q}_L \chi_R \bar{\chi}_R Q_L. \tag{5}$$

In analogy with Eq. (1), these interactions become critical and break EWS. We also allow the singlet mass terms $\bar{t}_R \chi_L$ and similarly $\bar{\chi}_L \psi_R$ and $\bar{\chi}_R \psi_L$. Finally we include the mass term $M^U \bar{\psi}_L \psi_R$, which provides the only connection between the left and right handed top quarks. Note that we have not included all possible singlet masses compatible with the gauge symmetries; effectively we have included one massive fermion ($\bar{\chi}_L \psi_R$) that couples to the t_L and another massive fermion ($\bar{\chi}_R \psi_L$) that couples to the t_R. The two massive fermions then have small mass mixings from the $\bar{\psi}_L \psi_R$ masses which are the only couplings between these sectors. Although this may seem somewhat *ad hoc* at this stage, the models of Secs. II and III naturally produce this mass structure. The important point here is that there is some singlet mass structure that can reproduce the SM fermion masses. The mass matrix takes the form

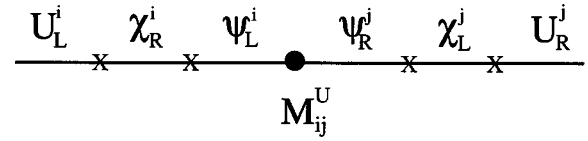

FIG. 1. The mixing that produces the up quark sector masses.

$$(\bar{Q}_L, \bar{\chi}_L, \bar{\psi}_L) \begin{pmatrix} 0 & m_1 & 0 \\ m_2 & 0 & m_3 \\ 0 & m_4 & M^U \end{pmatrix} \begin{pmatrix} t_R \\ \chi_R \\ \psi_R \end{pmatrix} \tag{6}$$

where m_1 is the EWS breaking mass. In general the matrix is complicated but the pattern of masses resulting can be seen by taking the decoupling limit. We imagine that the largest masses are m_3 and m_4 which bind the ψ and χ fermions into two heavy Dirac fermions. We may then treat the remaining masses as perturbations. The SM top quark mass results from the diagram in Fig. 1 and the mass is proportional to the singlet mass M^U.

This singlet sector mass structure is more readily convertible to include the other SM fermions. We will write the SM quarks as Q_L^i, U_R^i and D_R^i where i is a generation index. We may endow the fields χ and ψ with a generation index and the mass matrices $m_1 - m_4$ and M^U with structure in the generation space. A four fermion operator of the form of Eq. (5) for each left handed doublet will break EWS through a condensate for each flavor. We shall assume that these operators give rise to flavor universal VEVs since we expect the four fermion operators to arise from some flavor universal broken gauge interaction. The precise way in which this dynamics will be implemented depends on the choice of broken gauge symmetry and we will discuss two possibilities: a gauged *family* symmetry and the gauging of the complete *flavor* symmetry. The dynamical implementation of these models and their consequences are discussed in detail in the next two sections.

The important point of the present structure is that the masses of the light eigenstates are proportional to the singlet mass matrix M_{ij}^U. Thus, the suppression of the up quark mass relative to the top quark mass only requires us to have a relatively smaller mass element in the ψ mass matrix. The heaviest fermion masses (m_3, m_4) do not contribute to the generation of the structure in the SM masses and hence the upper cutoff of the theory can be left unchanged even when we incorporate the lightest fermion masses. There is therefore no increase in the degree of fine-tuning in the model if the up quark participates in EWS breaking. At least at this stage in the model building, the three generations of singlet fermions ψ reflect the same mass hierarchy problem present in the spectrum of SM fermions. The important point is that the origin of the symmetry breaking mass splittings has been deferred to this EWS singlet sector. It is not our intention to address their origin here, though one could imagine generating these masses dynamically, via additional interactions residing at even higher energies, such as some singlet sector extended techincolor model or tumbling gauge structure.

Extending this scenario of mass generation to the isospin $-1/2$ quark sector is straightforward. We introduce the electroweak singlets ω_L^i, ω_R^i, as well as ξ_L^i, ξ_R^i, all with the same hypercharge assignments as right-handed down quarks. A four fermion interaction $\overline{Q_L^i}\omega_R^i\overline{\omega_R^j}Q_L^j$ will dynamically generate the condensate $\langle\overline{Q_L}\omega_R\rangle$ contributing to EWS breaking. The singlets ξ have a mass matrix M_{ij}^D which gives the connection between the left and right handed down-quark sectors. M_{ij}^D generates the pattern of down quark masses through diagrams analogous to the one in Fig. 1.

In this naive model with a four fermion interaction and explicit singlet masses, a Cabibbo-Kabayashi-Maskawa (CKM) matrix may be easily incorporated. There are in total seven chiral family symmetries of the singlet fermions broken by six singlet mass matrices. These matrices cannot therefore in general be diagonalized simultaneously and the model contains mixing angles that will feed into the SM sector.

The electroweak singlets, χ and ψ in the up sector and ω and ξ in the down sector, have large masses relative to the weak scale (of the order of a few TeV) and may be considered to be decoupled. The dynamical breaking of electroweak symmetry by the condensates between the SM left-handed fermions and the χ_R, ω_R singlet fermions will generate a scalar sector of the low energy theory [3]. The precise form of the scalar sector is model dependent. We will discuss two examples of this sector in the following two sections where the origin of the dynamics is more concretely specified. However, it is possible to make a general point about the mass scale governing the masses of the scalar sector in relation to that of the top see-saw model. The important difference is that the EWS breaking VEV is flavr universal. In order to see this difference we notice that the Pagels-Stokar formula is given by

$$v^2=\frac{N_f}{4\pi^2}m_1^2\ln\left(\frac{\Lambda}{m_3}\right), \tag{7}$$

where N_f denotes the number of fermion flavors forming a condensate. For the top condensate model, $N_f=N_c$ and we recover Eq. (4). On the other hand, for example in a model where all the quarks participate equally in EWS breaking, $N_f=6N_c$. For instance, for $\ln\Lambda/m_3\simeq\mathcal{O}(1)$ one obtains $m_1\simeq 360$ GeV, whereas larger values of the ratio Λ/m_3 result in $m_1\simeq 250$ GeV or even lighter. This EWS breaking mass is approximately 3 times smaller than in the top seesaw model. If one naively computes the Higgs boson mass as in the large N approximation to the four fermion theory, the Higgs boson mass may be extracted from a bubble resummation and is found to be

$$m_h\simeq 2m_1, \tag{8}$$

implying Higgs boson masses in the range $m_h\simeq(500–700)$ GeV.

When the EWS breaking is driven by four fermion interactions it is easy to include the leptons (L_L^i, N_R^i and E_R^i) in the same pattern as the quarks. Additional EWS singlets are introduced analogous to χ^i, ψ^i, ω^i and ξ^i but with the SM interactions of the right handed leptons. Four fermion interactions analogous to those of Eq. (5) drive condensates involving L_L^i and the massive singlet sector. The standard model lepton masses come again from the mass matrices linking the singlet ψ fields. If the leptons participate equally in the EWS breaking condensates, then $N_f=6N_c+6$ in the Pagels-Stokar formula, and the Higgs boson masses are controlled by a even lighter scale ($m_h\simeq 400–650$ GeV).

One of the benefits of this class of models is that there are only very small deviations from the SM values of electroweak precision variables such as the S and T parameters. The S parameter essentially counts the number of electroweak *doublets*. The flavor universal model has only extra electroweak *singlet* fields and so makes no extra contribution. By taking the singlet mass terms (m_3,m_4) as very large one may decouple them from the low energy dynamics (perhaps at the expense of increased fine-tuning to generate v) or equivalently we may say that the physical heavy fermions have only a very small admixture of the electroweak doublets in them. In this limit the isospin breaking of the heavy fermions decouple from the T parameter [7], leaving just the contribution from the light SM fermions with the standard masses. As this limit is relaxed all the mass eigenstates, the heavy and the light, will contribute to T with by far the largest contribution from the top quark and its singlet partners. The calculation of T is very similar to that of the top seesaw model and the additional contributions to T are easily controlled to be of the order the experimental bounds for values of masses for the heaviest fermions of the order of (1–5) TeV.

In the next two sections we present specific examples of the flavor universal seesaw mechanism, where the broken gauge dynamics is made explicit. In one case it results from a broken gauged *family* symmetry, in the other from a larger broken gauged *flavor* symmetry. A crucial aspect of theories with gauged flavor symmetries is to avoid the tight constraints on FCNC. We will provide models that achieve this by maintaining a GIM mechanism above the electroweak scale and provide existence proofs that such gauge dynamics is possible. We also discuss the phenomenology of these models, considering the consequences of the broken flavor symmetry as well as the scalar and pseudo-scalar sector of the models.

III. DYNAMICS FROM BROKEN FAMILY SYMMETRY

Our first explicit example of a flavor universal EWS breaking model will assume that the family symmetry of the left handed SM fermions is gauged and broken at a scale of a few TeV. For simplicity we will restrict ourselves to the quark sector. The χ_R^i and ω_R^i fields are assumed to also transform under this gauge group. The broken gauged family interactions result in the four fermion interactions

$$\overline{Q_L^i}\chi_R^i\overline{\chi_R^j}Q_L^j+\overline{Q_L^i}\omega_R^i\overline{\omega_R^j}Q_L^j \tag{9}$$

where the family index is summed over. These interactions will generate the EWS breaking VEVs between the left handed quark multiplets and the singlet fermions.

The singlet sector masses are assumed to be in place to generate the quark mass matrices discussed in Sec. II.

The scalar sector of the theory is expected to be large. The quark condensates are breaking an $SU(6)_L \times SU(6)_R$ chiral symmetry of the quarks and χ_R, ω_R fields to the vector subgroup. To represent this as a Higgs model we must have a Higgs field transforming as a $(6,\overline{6})$ under the flavor group. There are thus 72 real scalars. As a result of the chiral symmetry breaking there are 35 NGBs of which three are absorbed to give masses to the W and Z gauge bosons. The remaining 32 pseudo-scalars are naively massless but acquire masses through SM gauge interactions. They correspond to a color-octet $SU(2)_L$ singlet and a color-octet $SU(2)_L$ triplet. Their masses, due to the strong interactions, can be computed to be [13]

$$m_\pi^2 = 3\alpha_s M^2, \tag{10}$$

where M is a high energy scale. For instance, for $M \simeq O(1)$ TeV, we obtain $m_\pi \simeq (500\text{–}600)$ GeV. They couple to SM quarks with couplings proportional to m_q/f_π (f_π is the PNGB decay constant $\sim v$).

As discussed in Sec. II the non-NGB scalars are expected to have masses in the (500–700) GeV range. One of these, a pseudo-scalar, is the would be NGB of the $U(1)_A$ symmetry; the remaining 36 scalars break down into the same SM gauge multiplets as the NGBs. We leave a more precise computation of the spectrum analogous to that of [3] for future investigation since it will be involved.

At the level of Eq. (9) the leptons may also be included in the interaction that breaks EWS. If the leptons participate in EWS breaking, then the flavor symmetry $SU(8) \times SU(8)$ is broken to the vector subgroup and there are 63 PNGBs of which again only 3 are absorbed. The additional 28 PNGBs consist of two electroweak triplets and two singlet leptoquarks with masses

$$m_\pi^2 = \frac{4}{3}\alpha_s \quad M^2, \tag{11}$$

that is $m_\pi \simeq (300\text{–}400)$ GeV; there is also a triplet and a singlet of color neutral PNGBs. These latter PNGBs do not receive masses from the SM gauge interactions at one loop and are potentially in conflict with experiment. This may be an indication that the leptons do not participate in EWS breaking but instead receive their mass by some radiative mechanism from the quark sector. In addition in this case there would be 65 massive scalars.

A. Dynamical model of the universal seesaw mechanism

In order to provide a more explicit, and renormalizable, model of the flavor universal seesaw mass pattern described in Sec. II we revive a model of top quark condensation by King (in turn based on the technicolor model of Georgi [9]). The model is shown in ''moose'' notation in Fig. 2. Recall that in moose notation a gauge group is represented by a circled number [N for $SU(N)$] and left-handed fermions transforming in the (antifundamental) fundamental representation of that group as lines with (outward) inward pointing arrows. Global symmetry groups are un-circled numbers. For simplicity we only consider the quark sector of the theory until the end of this section.

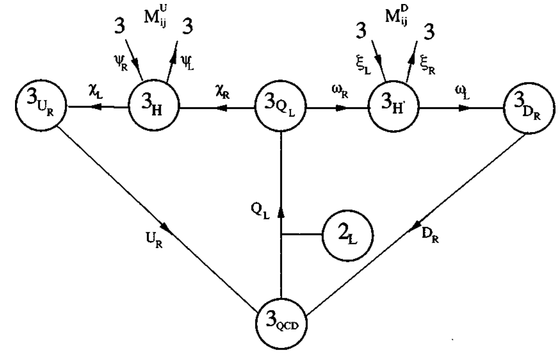

FIG. 2. A dynamical model of the universal seesaw mechanism for the quark sector displayed in moose notation.

In Ref. [8] the dynamics was assumed to occur in the following fashion. We concentrate on the isospin $+1/2$ sector for ease of discussion. The $SU(3)_H$ ''hypercolor'' gauge group under which χ and ψ transform was assumed to get strong, generating the condensates $\langle \overline{\chi_L} \psi_R \rangle$ and $\langle \overline{\chi_R} \psi_L \rangle$ and breaking the two chiral family symmetry $SU(3)$ groups into global groups. The formation of these condensates is not the preferred breaking pattern when the family symmetry groups are weakly coupled. However, there is some evidence that the effective potential favors this pattern when the family symmetry groups are strongly coupled at the symmetry breaking scale. This hypothesis was first put forward by Georgi [9], elaborated on in Ref. [14] and disputed in Ref. [15]. We shall assume that this pattern of condensation is correct. In the model of King these condensates were the only ones to form at this scale and served the purpose of breaking the family gauge groups, leaving strongly coupled ''flavorons.'' The flavorons of the two chiral family groups mixed through the $\overline{\psi_L}\psi_R$ mass matrix and the result was four fermion interactions between the isospin $+1/2$ quark sector with the strongest channel proportional to the largest eigenvalue of the ψ mass matrix. This strongest interaction was assumed to generate a top quark condensate $\langle \overline{t_L} t_R \rangle$.

We now propose an alternative possibility. Since the family symmetry groups are strongly interacting and close to their chiral symmetry breaking scale when the hypercolor group breaks its chiral symmetries, it is not clear that the χ and ψ fields should be integrated out of the dynamics of the family gauge groups. It is very plausible that the condensates $\langle \overline{U^i_R} \chi_L \rangle$ and $\langle \overline{U^i_L} \chi^i_R \rangle$ form as well. If we assume this, then the strong dynamics immediately produces the pattern of masses needed for the flavor universal seesaw model. Precisely how large the three different condensates are is a sensitive matter of the three gauge group interrelated strong interactions. We shall assume that by appropriate choices of couplings at the hypercolor group's strong scale the desired values can be obtained. There is of course, as in all top condensate models envisaged to date, a serious problem of fine-tuning in the assumption that many gauge groups become strongly interacting at almost an identical scale. We offer no excuses (though it might be more natural if the gauge dynamics is walking in nature) but put forward this model as a simple existence proof of dynamics that might give rise to the required effective theory.

The existence of these additional condensates also breaks the $SU(3)$ color group and the hypercolor group to their vector subgroup that then plays the role of the low energy QCD group. There are thus flavor universal colorons analogous to those in the model of Ref. [10]. The NGBs associated with the condensate $\langle \overline{U}^i_R \chi^i_L \rangle$ are absorbed in the breaking of these two groups to the vector subgroup. The NGBs from the condensate $\langle \overline{U}^i_L (\chi^i_R + \omega^i_R) \rangle$ are those already discussed above. Finally there are the 36 NGBs associated with the $\langle \chi\psi \rangle$ condensates of which 16 are absorbed to give masses to the two sets of family symmetry group gauge bosons. The remaining NGBs will acquire masses from the explicit mass matrix M^U_{ij} and the family gauge interactions that explicitly break the chiral symmetry group. However, the mass matrix is off diagonal in the basis of the condensate and the gauge interactions chiral; so these symmetry breakings will only contribute to the PNGB masses at second order. These EWS singlet PNGBs may therefore only have masses of the order of a few GeV. We will live with their existence since they only occur in the singlet sector of this model, which is in any case only intended to illustrate the concept of a flavor universal seesaw mechanism. We do not believe they are necessarily a robust signature of this class of models.

The bottom quark sector of the model behaves similarly to the top quark sector and the relative mass differences in the final spectrum are trivially the result of different mass matrices included by hand for the ψ and ξ singlet fermions of the two sectors. Including mixing angles is non-trivial [9]. They will result if the singlet fermion mass matrices in the up and down sectors, M^U_{ij} and M^D_{ij}, are not simultaneously diagonal. As shown in Fig. 2 the model has sufficient global symmetry to simultaneously diagonalize both mass matrices and there will be no mixing angles. An attempt to remedy this problem by including yet higher energy scale dynamics can be found in Ref. [9]. We wish to defer questions of such high scale dynamics and simply note that there does not appear to be any problem, in principle, preventing the generation of mixing angles from the mass matrices of the singlet sector.

The model has flavor symmetries gauged at a relatively low scale, $\mathcal{O}(1)$ TeV, which might naively be expected to generate FCNC. However, as promoted in Ref. [9], the model has the same chiral flavor symmetries as the SM even above the weak scale (that is $[SU(3)\times U(1)]^3$) and hence a GIM mechanism that allows the neutral current to be written in diagonal form. FCNC effects will be generated with the inclusion of the generic mass matrices $M^{U/D}ij$ but the leading effects are through the quark masses themselves, and therefore simply correspond to the SM contributions. The largest additional contributions are from mass mixings of the left and right family gauge groups that lead to four fermion operators between the up- (down-) type quarks mixing through $M^D(M^U)$ which rotations on those fields cannot diagonalize. The leading such term contributing to $\Delta S=2$ processes is for example

$$\frac{1}{\Lambda_H^6}\,\overline{Q}^i_L (M^U M^{U\dagger})_{ij} Q^j_L \overline{Q}^k_L (M^U M^{U\dagger})_{kl} Q^l_L . \qquad (12)$$

In addition to the suppression by the gauge boson mass these contributions are suppressed by a factor of the SM mass mixing matrix element over 50 GeV to the fourth power. They are therefore much smaller than the SM contributions.

The inclusion of leptons in the model is not so straightforward. Naively one would simply repeat the model of the quark sector with the SM leptons transforming under the same family symmetry groups. However, this would introduce unacceptably large lepton flavor violating effects (e.g. $K^0 \to e\mu$) since two multiplets would now transform under a single family symmetry group only broken at the few TeV scale. As an alternative, one could simply replicate the whole moose model for the lepton sector. Additional χ and ω fields transforming under their own hypercolor groups would need to be added and the lepton mass matrices would come out proportional to the mass matrix put in between the new ω fields. The only problem is that to cancel anomalies the hypercolor group must be a $U(1)$ group which is not asymptotically free. Rather than make the model more baroque to force the leptons in, we will leave the model at this point since the similar model of the next section succeeds more naturally in including the leptons, although not without similar problems.

B. Phenomenological considerations

In this section we discuss the phenomenological constraints on the family symmetry generated flavor universal seesaw model as well as its potential signatures at future experiments. Here we concentrate on the model giving quark masses and EWS breaking, leaving the discussion of lepton masses for the next section. The model described in the previous section contains several new states. There are the fermion singlets, $\chi_{L,R}$ and $\psi_{L,R}$ connected to the up-quark sector and $\omega_{L,R}$ and $\xi_{L,R}$ to the down-quark sector. As mentioned in Sec. III, these acquire large masses as a result of the condensates formed among them, induced by the strong up and down hypercolor groups. This, together with the fact that they are electroweak singlets, implies that they are not relevant to the low energy phenomenology. They are, at this stage, simply conduits to communicate between the left- and right-handed quark sectors and to mix with the left-handed quark doublets and generate EWS breaking at the right scale.

1. Flavorons

The breaking of the three gauged family groups $SU(3)_L$, $SU(3)_{U_R}$ and $SU(3)_{D_R}$ by the formation of the necessary singlet fermion condensates implies the existence of three new sets of massive gauge bosons. These are family octets coupling separately to left-handed quark doublets and right-handed up and right-handed down quarks. The masses of these ''flavorons'' are expected to be similar and are determined by the scale where the hypercolor groups break the family symmetries.

The couplings of the flavorons to the SM fermions are given by

$$\mathcal{L}_f = -g_L L_\mu \overline{Q}^i_L \gamma^\mu t^A_{ij} Q^j_L - g_U R^U_\mu \overline{U}^i_R \gamma^\mu r^A_{ij} U^j_R - g_D R^D_\mu \overline{D}^i_R \gamma^\mu s^A_{ij} D^j_R, \tag{13}$$

where L_μ, R^U_μ and R^D_μ are the flavorons corresponding to the gauge family groups $SU(3)_L$, $SU(3)_{U_R}$ and $SU(3)_{D_R}$ respectively, and t^A, r^A and s^A are their respective generators, with $A=1,\ldots,8$. They simply are $t^A_{ij}=r^A_{ij}=s^A_{ij}=\lambda^A_{ij}/2$, with λ^A the generators of $SU(3)$. The fact that flavorons are flavor-octet gauge bosons seems to indicate the existence of non-diagonal vertices. However, these do not result in FCNC processes due to the presence of the remnant global $SU(3)$ family symmetries after the breaking of the gauged ones. These approximate global symmetries allow us to implement the GIM mechanism along the lines of the models of Ref. [9] and are only broken by fermion mass terms.

The flavorons acquire masses at the scale Λ_f where the hypercolour groups break the family symmetry. Thus, they are expected to be in the TeV range. On the other hand, their couplings to fermions must be strong enough to generate the $\langle \overline{Q}_L \chi_R \rangle$ and $\langle \overline{Q}_L \omega_R \rangle$ condensates that break the EWS, as well as the $\langle \overline{U}_R \chi_L \rangle$ and $\langle \overline{D}_R \omega_L \rangle$ necessary to obtain the SM quark masses. We compute the criticality condition in the Nambu–Jona-Lasinio (NJL) approximation. Defining $\kappa_a \equiv g_a^2/4\pi$, with $a=L,U,R$, the couplings must satisfy

$$\kappa_a \geq 2\pi \left(\frac{N_g}{N_g^2-1} \right), \tag{14}$$

where N_g refers to the number of generations. The condition, Eq. (14), has important phenomenological consequences. For instance, it gives a lower bound for the production cross section, for a fixed flavoron mass. At the same time it implies that the flavoron widths are rather large. For instance the flavorons coupled to the left-handed quarks have a width given by

$$\Gamma_f \simeq \frac{\kappa_L}{6} M_F, \tag{15}$$

which implies that the minimum width, given by the critical value of κ_L, is approximately 40% of its mass, making it difficult for them to be detected as clear mass bumps in $p\bar{p}$ collisions at the Tevatron. The situation is somewhat better for the flavorons coupled to the right-handed up and down quarks. Their minimum widths are about half that of Eq. (15). On the other hand and for all cases, the large couplings in Eq. (14) ensure large excesses in hadronic production of all flavors. The fact that flavorons are chirally coupled may produce a very distinct signal in high transverse momentum jets. At energies below the flavoron mass these excesses mimic the behavior of contact terms with specific chirality. The limits from the Tevatron run I data as well as the reach of run II are currently under study.

The presence of strongly coupled gauge bosons may be constrained by their contributions to the T parameter. As pointed out in Ref. [4], this is so even when the interaction itself is isospin conserving. The exchange of the new gauge bosons gives a two loop radiative correction to the isospin violating contributions to the ρ parameter coming from one-loop diagrams involving the SM fermions. We first notice that these types of contributions to T are induced only by the gauge bosons that couple to left-handed fermions. Thus, only one of the three flavorons must be considered. The relevant contribution to the T parameter is given by

$$T = \frac{4\pi}{s^2\theta_W c^2\theta_W M_Z^2} \left(\frac{\Pi_{LL}(m_t,m_b)}{2} - \frac{\Pi_{LL}(m_t,m_t)}{4} \right), \tag{16}$$

where $\Pi_{LL}(m_t,m_b)$ and $\Pi_{LL}(m_t,m_b)$ refer to the left-handed vacuum polarizations involving one top and one bottom quark, and two top quarks in the loop respectively, and they are evaluated at zero momentum transfer. Following [4] we will approximate the calculation of the two-loop diagram by a product of two one-loop diagrams obtained after shrinking the flavoron propagator. Then, the vacuum polarizations are

$$\Pi(m_t,m_b) = \frac{\Pi(m_t,m_t)}{4} \simeq -\frac{1}{16\pi^3} m_t^4 \left(\log\frac{\Lambda_f}{m_t} \right)^2 \frac{\kappa_L}{M_F^2}, \tag{17}$$

where the vacuum polarizations were approximated by their divergent behavior given by the leading logarithm. This results in

$$T = \frac{m_t^4}{s^2\theta_W c^2\theta_W M_Z^2} \frac{1}{8\pi^2} \left(\log\frac{\Lambda_f}{m_t} \right)^2 \frac{\kappa_L}{M_F^2}. \tag{18}$$

This constraint is compatible with Λ_f being of the order of the TeV scale. For instance, if we take $\log(\Lambda_f/m_3)\simeq 1$, and we assume $m_3 \simeq 1$ TeV, then the induced T parameter is

$$T \simeq 0.06 \frac{\kappa_L}{M_F^2}. \tag{19}$$

The current determination of T from electroweak precision measurements gives [16] $T=(-0.11\pm 0.16)$, assuming a 300 GeV Higgs boson mass. Thus, we see from Eq. (14) and Eq. (19) that flavoron masses in the 1 TeV range and above are not in contradiction with the electroweak precision data.

2. Colorons

The flavorons induce four-fermion interactions between quarks and singlet fermions which are supercritical and result in EWS breaking as well as in the breaking of the hypercolor groups in such a way that $SU(3)_H \times SU(3)_{H'} \times SU(3)_c$

breaks down to ordinary QCD interactions. This breaking pattern leaves two sets of color-octet massive gauge bosons, the ''colorons,'' each of them coupling separately to the up and down quarks.

The two sets of colorons in this model present some similarities with the massive color-octet gauge boson in the model of Ref. [10]. As with the flavorons, they will give contributions to the renormalization of SM ρ parameter. The resulting bounds are similar to those discussed for the flavorons and are in the TeV range.[2] Moreover and similarly to the case of the flavorons discussed above, the coloron couplings must be supercritical in order for the hypercolour groups to generate mass terms in the singlet sector. This constraint, here once again, will imply large production cross sections but also large widths. The width of either of the color octets is given by $\kappa_C \; M_C$, with κ_C, the coupling of the corresponding coloron, reflecting the embedding of QCD in the hypercolor group and above the critical value. The main difference with the flavoron case is that colorons interfere with the QCD gluon-mediated processes. The consequences of this interference at hadron colliders were first studied in Ref. [17] in the context of top-color models, and later investigated in Ref. [10] for the case of flavor universal color-octet gauge bosons. The phenomenology of the coloron sector of the model is very similar to the one described in this latter work.

3. Scalars and PNGB

Finally, in the model of Fig. 2 the breaking of the $SU(6)_L \times SU(6)_R$ chiral symmetry under which Q_L, χ_R and ω_R transform, by a four fermion interaction, implies the existence of a $(6,\overline{6})$ Higgs boson. The 72 real degrees of freedom split into 3 NGBs absorbed by the W and Z bosons, 32 color-octet PNGBs with masses estimated above to be in the few hundred GeV range, one PNGB acquiring a large mass due to the $U(1)$ anomaly, and 36 scalars in the Higgs sector with masses estimated to lie in the range (500–700) GeV.

These states couple to the SM fermions through Yukawa interactions of order m_f/f_π. In fact they are precisely the lightest states found in a one family technicolor model without techni-leptons, and are subject to the similar low energy constraints. The most significant bounds on this scalar sector come from the measurements of $\Gamma(Z \to b\overline{b})$ and the $b \to s\gamma$ branching ratio. The dominant contributions correspond to the color-octet scalars and pseudo-scalars. For instance, the effect of considering one of these color-octets in the $b \to s\gamma$ transition, when combined with the 3σ interval from the most recent CLEO measurement of the inclusive rate [18], $\mathrm{Br}(b \to s\gamma) = (3.15 \pm 0.35 \pm 0.32 \pm 0.26) \times 10^{-4}$, translates into a lower bound of 440 GeV on the mass of the charged states. The effects of adding all the scalar and pseudo-scalar states may tighten the bounds, depending on the scalar masses. However, large cancellations are also possible among the various contributions and the SM $b \to s\gamma$ amplitude. Additional information from $b \to s l^+ l^-$ decays, to be available in the near future from the next generation of B physics experiments, will elucidate this possibility. On the other hand, the R_b bounds are not affected by the these cancellations. For instance, the 3σ constraint on the color-octet states (assuming equal masses for scalars and pseudo-scalars) translates into a mass limit of 700 GeV. However, if the central value of R_b is taken to be the SM one, the mass constraint is considerably smaller, although still a few hundred GeV.

[2] The essential difference between the two cases resides in the fact that the same colorons couple to left and right handed fermions. As a result, the effect in Eq. (19) is larger by a factor of 3 for colorons.

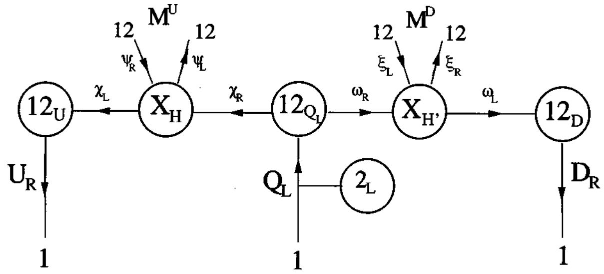

FIG. 3. A flawed dynamical model of the universal seesaw mechanism resulting from broken flavor symmetry displayed in moose notation.

IV. DYNAMICS FROM BROKEN FLAVOR SYMMETRY

There is an alternative possible flavor structure for the four fermion interactions driving EWS breaking that allow the leptons to condense and hence acquire masses but without giving rise to the light PNGB found in Secs. II and III above. That is to allow the index i on Q_L,U_R and D_R to run over the full possible set of fermion flavors. In other words, in this model the index i runs over the full $SU(12)$ flavor symmetry of the SM where the 12 flavors are three families of three colors of quarks and the three families of leptons. We introduce additional χ^i (ω^i) and ψ^i (ξ^i) electroweak singlets with a similar pattern of singlet masses to those of the model of Sec. II. The color and hypercharge interactions are weakly gauged subgroups of the $SU(12)$ flavor symmetry of the SM particles and the electroweak singlet sectors. The four fermion interaction

$$\overline{Q^i_L}\chi^i_R\overline{\chi^j_R}Q^j_L + \overline{Q^i_L}\omega^i_R\overline{\omega^j_R}Q^j_L \qquad (20)$$

drives electroweak symmetry breaking through the condensate $\langle \overline{Q^i_L}(\chi^i_R + \omega^i) \rangle$. The $SU(2) \times SU(2)$ chiral symmetry of the fermion doublets is broken to the vector subgroup producing three Goldstone bosons that are all absorbed by the W and Z bosons. In this case there are two charged Higgs bosons, two neutral Higgs and a massive pseudo-scalar, as in a standard two Higgs doublet model, with masses of the order of (400–650) GeV and no PNGBs, which makes this model phenomenologically more appealing than the one in the previous section.

It is natural to promote this idea to a moose model based on that of Sec. III. One might propose the moose model of Fig. 3, a variant of the technicolor model of Ref. [11]. We expect again that the dynamics is such that the $SU(X)$ hypercolor groups become strong and break their chiral symmetries in the few TeV range. The chiral flavor groups are also assumed to be strong at this scale. Amid these dynamics the condensates $\langle Q^i_L(\chi^i_R+\omega^i_R)\rangle$, $\langle U^i_R\chi^i_L\rangle$ and $\langle D^i_R\omega^i_L\rangle$ are assumed to form. Mass matrices put in by hand for the ψ and ξ fermions provide the left-right connecting structure necessary to produce the SM fermion masses. This is just the universal seesaw mass scheme. The SM gauge interactions result as the vector subgroup of the gauged $SU(12)$ chiral flavor symmetries and any weakly gauged subgroup of the ψ and ξ fermions global $SU(12)$ symmetry groups. We do not show these weakly gauged subgroups in the figure for ease of display. The electroweak symmetry breaking VEVs do not generate any PNGBs since the 3 NGBs arising from the breaking of the $SU(2)_L\times SU(2)_R$ chiral group are absorbed by the W and the Z gauge bosons. However, the hypercolor dynamics will generate 575 NGBs, corresponding to the breaking of the $SU(24)_L\times SU(24)_R$ chiral symmetry of the χ's and the ψ's. Of these, only 286 are absorbed by the broken flavor symmetry groups. The remainder are, just as in the quark moose, potentially light although they are singlets under the SM interactions. As in that model, these light states are a result of the particular implementation of the singlet sector and not an automatic consequence of the flavor universal models. The model again has a GIM mechanism that suppresses FCNC.

A more serious problem with the present model resides in the fact that anomaly cancellation for the gauged $SU(12)$ groups requires $X=1$. Attempts to make these groups non-Abelian by the inclusion of new EWS singlet fermions transforming under the $SU(12)$ groups give rise to additional PNGBs, which may be unacceptably light (see Ref. [11] for attempts in these directions).

We present Fig 3, although it is unsatisfactory, because it helps to visualize the dynamics we envisage might be behind the model presented first at the four fermion level. A flavor universal EWS breaking scenario is possible without PNGBs and will most likely be associated with a large, broken, chiral, gauged flavor symmetry of at least the left handed SM fermions. The hypercolour dynamics is not essential to the basic idea and may at least temporarily be replaced by Higgs bosons in the absence of model building ingenuity.

The broken flavor symmetry model is much cleaner at low energies than the model of Sec. III. There are no unabsorbed NGBs associated with EWS breaking. The dynamics is expected to give rise to four scalar Higgs bosons and a pseudo-scalar each with masses (400–650) GeV which correspond to a standard two Higgs doublet model with the ratio of Higgs VEV, $\tan\beta=1$.

The massive flavorons are similar to those of the previous model except that they now also mediate interactions involving leptons. The GIM mechanism of the model suppresses FCNC. One important modification with respect to the model of the previous section is the criticality condition. The couplings $\kappa_{L,R}$ obey conditions similar to Eq. (14) but with N_g replaced by $N=12$. As a result the necessary couplings are considerably smaller than in the previous cases, which has important consequences in the phenomenology. For instance, the minimal flavoron contributions to the T parameter in Eq. (19) can now be smaller by roughly a factor of 4, relaxing the mass bound by a factor of 2. On the other hand, somewhat more stringent limits come from neutral current processes involving effective contact terms of the form $ll\overline{qq}$. We concentrate on $l=e$ where most of the experimental information is. The effective $eeqq$ interactions are parametrized by the standard expression [19]

$$\mathcal{L}_{\rm NC}=\sum_q \{\eta^{eq}_{LL}(\overline{e_L}\gamma_\mu e_L)(\overline{q_L}\gamma^\mu q_L)+\eta^{eq}_{RR}(\overline{e_R}\gamma_\mu e_R)\times(\overline{q_R}\gamma^\mu q_R)+\eta^{eq}_{LR}(\overline{e_L}\gamma_\mu e_L)(\overline{q_R}\gamma^\mu q_R)+\eta^{eq}_{RL}(\overline{e_R}\gamma_\mu e_R)(\overline{q_L}\gamma^\mu q_L)\}. \tag{21}$$

The exchange of the $SU(12)$ flavorons generates the first two terms, whereas the LR and RL terms arise only as a consequence of one loop mixing of the L and R gauge bosons, which is suppressed by two powers of the masses, relative to the LL and RR coefficients. The LL and RR coefficients take the form

$$\eta^{eq}_{LL}=\frac{\pi\kappa_L}{6M^2_{F_L}}, \tag{22}$$

$$\eta^{ed}_{RR}=\frac{\pi\kappa_R}{6M^2_{F_R}}, \tag{23}$$

where $q=u,d$. The $SU(2)_L$ relations naturally resulting in this model imply that the RR interactions of the u quark involve neutrinos instead of charged leptons. This mismatch results in potentially dangerous contributions to which atomic parity violation (APV) as well as neutrino scattering experiments are especially sensitive. The most constraining bounds come from the APV experiments. The resulting contribution to the atomic weak charge is given by [20]

$$\Delta C_q=\frac{v^2}{2}(\eta^{eq}_{RR}-\eta^{eq}_{LL}). \tag{24}$$

Thus, assuming that the couplings and masses of the left and right handed sectors are similar, the effect largely cancels in the $eedd$ interactions. This is not the case for $eeuu$. From the most recent measurements [20] we obtain the 3σ bound

$$\eta^{eu}_{LL}<0.22 \tag{25}$$

which, for the critical value of the flavoron coupling, translates into $M_F>1.2$ TeV. Looser bounds on M_F are obtained from the νN scattering experiments, as well as from the DESY collider HERA, LEPII and Drell-Yan processes at the Tevatron. Thus, a scale of a few TeV for the dynamics of this model is not in contradiction with experimental observations.

V. CONCLUSIONS

We have presented model building ideas that extend the top condensate seesaw model of Ref. [7] to include realistic mass generation for all the SM fermions. In these models all the electroweak doublets of the SM play an equal role in breaking EWS, with their mass differences resulting from different mass mixings in a heavy EWS singlet fermion sector. This effectively separates the questions of EWS breaking and the origin of fermion masses, deferring the latter to higher energy scales. As a result, the origin of the dynamics of fermion masses is removed as a source of contamination of the precision electroweak variables S and T. The mass mixings of the EWS singlet fermions, which are the seed for the standard model fermion masses, may be generated at much higher scales than the weak scale without inducing further fine-tuning in the production of the weak scale.

We presented two classes of models, based on broken gauged family symmetry in one case and broken gauged flavor symmetry in the other. In both cases the resulting composite Higgs sector is relatively light, with masses in the range $m_h = (400–700)$ GeV. In gauged family models, the breaking of the large chiral symmetries in the SM fermion sector leads to the presence of a large number of scalars and pseudo-scalars with masses in this range, resulting in constraints from the measurements of R_b and FCNC processes. On the other hand, the scalar sector of the broken flavor symmetry models is more economical, with no PNGBs, and corresponds to the two Higgs doublet model with $\tan\beta = 1$.

Both scenarios involve extensions of the gauge sector and imply the existence of massive gauge bosons associated with the breaking of either the family or flavor symmetries, the flavorons. These are strongly coupled to the SM fermions and thus have a very rich phenomenology at present and future high energy colliders. We have seen that the mass scale of the flavorons can be as low as a few TeV and still satisfy all existing phenomenological bounds. Therefore, the experiments to take place at the Fermilab Tevatron Collider, as well as the CERN Large Hadron Collider (LHC), will be sensitive to a variety of signals associated with these states, as well as with the colorons already present in other models of dynamical EWS breaking [21].

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