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Electroweak symmetry breaking from a holographic fourth generation

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Abstract: We consider a model with four generations of standard model fermions propagating in an extra dimension with an AdS background metric. We show that if the zero modes of the fourth generation are highly localized towards the infrared brane, it is possible to break the electroweak symmetry via their condensation, partly driven by their interactions with the Kaluza-Klein excitations of the gauge bosons, as well as by the presence of bulk higher-dimensional operators. This dynamical mechanism results in a composite Higgs, which is highly localized and generally heavy. The localization of fermions in the five-dimensional bulk naturally leads to the standard model Yukawa couplings via the action of the bulk higher-dimensional operators, which are suppressed by the Planck scale. We obtain the spectrum of the model and explore some of its phenomenological consequences, both for electroweak precision constraints as well as at the Large Hadron Collider.

Keywords: Field Theories in Higher Dimensions, Beyond Standard Model, Technicolor and Composite Models.

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Contents

1. Introduction

The standard model (SM) is an extremely successful description of the electroweak interactions. However, the instability of the weak scale under radiative corrections leads us to believe that there should be physics beyond the SM at an energy scale not far beyond the TeV. The origin of electroweak symmetry breaking (EWSB) as well as of fermion masses, might be associated with this new dynamics. For instance, the fact that the top mass is of the order of the weak scale suggests that its origin might be associated with EWSB. This was first proposed in ref. [1], where the condensation of the top quark leads to its dynamical mass and the breaking of the electroweak symmetry. The new dynamics responsible for the condensation cannot be far above the weak scale in order to avoid fine-tuning. However, in order for the new dynamics to occur close to the TeV scale, the top-quark's dynamical mass would have to be considerably larger than what it actually is, about $(600-700)\ GeV$. Conversely, in order to obtain the experimentally observed m_t, the new dynamics has to reside at a scale of about 10^{15} GeV or so. Already in ref. [1] it was pointed out that the condensation of a fourth generation, with TeV dynamics, could be the solution to this problem. In this context, however, it is not simple to arrange the dynamics so that only the fourth generation condenses, as well as to explain the rest of the fermion masses [2]. The strong dynamics appears unnaturally selective of the different fermion generations.

Non-trivial strong dynamics in 4D can be described by a weakly-coupled theory in 5D through holography [3]. Of particular interest is the case in which the metric of the extra

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dimension is Anti-de-Sitter (AdS). This choice leads to the possibility of large separation of energy scales and has been proposed to solve the hierarchy problem [4]. In a compact extra dimension with AdS metric, and the fundamental scale being the Planck mass M_P, it is possible to generate the TeV scale at a distance πR from the origin, as long as $kR \sim (10-12)$, where k is the AdS curvature $k \sim M_P$. This type of setup not only can explain the hierarchy between M_P and the weak scale. If fermions are allowed in the bulk [5−8], their localization in the extra dimension, determined by their $O(M_P)$ bulk masses, results in exponentially separated overlaps with the TeV scale, which would explain the fermion mass hierarchy naturally. In this scenario the Higgs field must remain highly localized close to the TeV brane in order not to receive quadratic divergences to its mass above the TeV scale. In principle, this localization should have a dynamical origin. At the moment there is only one dynamical mechanism localizing the Higgs field. This naturally occurs in models where the Higgs is obtained from a higher dimensional component of a bulk gauge field [9, 10]. In this case, the Higgs corresponds to the zero mode of the part of the bulk gauge field related to the broken generators. This essentially means that this Higgs is a pseudo-Nambu-Goldstone boson. Alternatively, Higgsless scenarios [11] have been proposed in AdS_5, where the electroweak symmetry is broken by boundary conditions. Finally, it is possible to interpolate between these two pictures [12] by having a bulk Higgs with a TeV-localized vacuum expectation value (VEV).

In this paper, we consider four SM generations propagating in an AdS_5 bulk. The bulk masses of the fourth-generation are chosen so as to localize its zero-modes towards the TeV brane. This in turn induces strong couplings of the fourth-generation to the Kaluza-Klein (KK) excitations of the gauge bosons, particularly of the fourth-generation quarks with the KK gluons. Also, the inevitable presence of bulk higher-dimensional operators induces additional zero-mode four-fermion operators. The effectively induced four-fermion interactions can be super-critical, breaking chiral symmetry and the electroweak symmetry. In this realization of the fourth-generation condensation, the obtained dynamical fermion mass is approximately $(600-700)$ GeV, for KK gauge masses in the few TeV region. In the simplest realization, with only one fourth-generation zero-mode quark condensing (e.g. the up-type), the effective theory at energies below the KK mass scale presents a spectrum containing only one composite scalar doublet corresponding to the Higgs field. As long as the four-fermion interactions induced by the KK excitations are super-critical in the condensing channel, the Higgs acquires a VEV, and the electroweak symmetry is broken, giving masses to the $W^{\pm}$ and Z^0. The TeV localization of the Higgs field is a direct result of the localization of its constituents. We find typically a heavy Higgs, as is to be expected due to its highly TeV-localized wave-function, about 900 GeV for a few TeV KK masses.

Bulk four-fermion operators, suppressed by the Planck scale, will be responsible for fermion masses. In particular, four-fermion operators involving the condensing fourth generation quarks will result in fermion masses. Just as in any bulk Randall-Sundrum (RS) model with a TeV-localized Higgs, fermion zero modes with large overlap with the TeV brane will be heavier (e.g. the fourth generation, the top quark), whereas Planck-brane localized fermions will have suppressed coefficients in the four-dimensional effective operators resulting from the higher dimensional 5D operators. Thus, the model maintains the natural

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generation of the fermion mass hierarchy, a very compelling feature of bulk RS models.

The model we present here is a realization of the flavor-dependent strong dynamics necessary in a fourth-generation condensation scenario, in the context of the AdS/CFT correspondence. It also provides an alternative way to localize the Higgs field close to the TeV brane in RS models, other than the one proposed in refs. [9, 10]. Also, unlike in the model of ref. [12], the Higgs VEV and its localization are not free parameters, but are fixed by the dynamics of the fourth-generation in the bulk. The general idea allows for various choices, from the number of condensing fermions, to the presence of a right-handed neutrino zero-mode, and generally the choice of fermion representations under the bulk gauge group. We will try to be as definite and simple as possible, leaving alternatives for further work.

The plan for the paper is as follows: in the next section we present the model and show how electroweak symmetry and fermion masses arise in it. In section 3, the low energy effective theory for the zero modes and the Higgs is built. We compute the masses of the fourth-generation fermions and the Higgs making use of renormalization group methods. In section 4 we consider the electroweak precision constraints on the model, and its main phenomenological features, especially at colliders, are discussed in section 5. We conclude in section 6.

2. The model

2.1 The five-dimensional setup

We consider a theory with one compact extra dimension where the metric is given by [4]

$$ds^2 = e^{-2\,k\,y}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2\,, \tag{2.1}$$

and $k \sim M_P$ is the AdS curvature. The orbifold compactification S_1/Z_2 results in a slice of AdS in the interval $[0, \pi R]$, with R the compactification radius. In order for the weak scale to arise at the brane in $y = \pi R$, we need $k \sim 11$. All fermions propagate in the 5D bulk. These include the standard three generations, as well as a complete fourth generation. The boundary conditions are such that the zero-mode spectrum reproduces that of the SM fermions, with the addition of one extra SM generation. The gauge symmetry in the bulk cannot be just the SM, since the custodial breaking contributions from the $\mathrm{U}(1)_Y$ KK modes would result in unacceptably large isospin violation. Instead, we consider that the bulk gauge theory is [13] $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_X$, where the boundary conditions in the UV lead to $\mathrm{SU}(2)_R \times \mathrm{U}(1)_X \to \mathrm{U}(1)_Y$. Additional symmetries may be imposed in order to protect the $Z\bar{b}b$ coupling from large corrections [14]. When this is the case, third generation fermions have to be in specific representations of the gauge group. For instance, left-handed quark doublet must be a $(\mathbf{2},\mathbf{2})_{\mathbf{2/3}}$ under $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_X$. On the other hand the field with t_R as its zero mode, can be in either $(\mathbf{1},\mathbf{1})_{\mathbf{2/3}}$ or $(\mathbf{3},\mathbf{1})_{\mathbf{2/3}} \oplus (\mathbf{1},\mathbf{3})_{\mathbf{2/3}}$. We will take similar representations for the fourth generation. Regarding leptons, we will assume at a minimum the presence of a fourth generation lepton bi-doublet, and a singlet with a charged right-handed zero-mode E_R. Also, if we assume that the fourth-generation neutrino has a large Dirac mass, there should be an additional bulk field with a right-handed neutrino zero-mode N_R. Since in this paper we are mainly concerned with the fourth-generation zero modes, we will not need to make a choice of bulk representation, whenever such choice is possible.

Bulk fermion masses are naturally of the order of the AdS curvature k, such that

$$M_f = c_f\, k\,, \tag{2.2}$$

with $c_f \sim O(1)$. They determine the localization of fermion zero modes in the bulk. The localization of the fourth-generation zero-modes very close to the TeV brane results in strong interactions with the gauge boson KK modes. The couplings of fermions to KK gauge bosons are generically determined from the expression for the 5D coupling

$$g_5 \int d^4x \int_0^{\pi R} dy \sqrt{g}\bar{\Psi}(x,y)\, e^{ky}\, \gamma^\mu T^a \Psi(x,y) A_\mu^a(x,y)\ . \tag{2.3}$$

where the factor of e^{ky} comes from the vierbein, g_5 is the 5D gauge coupling, and the T^a are the generators of the gauge symmetry. Expanding $A_\mu(x,y)$ and $\Psi(x,y)$ in their KK modes as

$$A_\mu(x,y) = \frac{1}{\sqrt{\pi R}} \sum_n \chi_n(y) A_\mu^{(n)}(x)\,, \tag{2.4}$$

and

$$\Psi(x,y) = \frac{1}{\sqrt{\pi R}} \sum_n e^{2ky}\, f_n(y) \psi^{(n)}(x)\,, \tag{2.5}$$

and integrating over the compact dimension we obtain the coupling of the ith fermion KK mode to the n-th KK mode of the gauge boson:

$$g_{in} = \frac{g}{\pi R} \int_0^{\pi R} dy\, e^{ky}\, |f_i(y)|^2\, \chi_n(y)\,, \tag{2.6}$$

where we have defined the 4D gauge coupling by $g = g_5\,/\sqrt{\pi R}$. Here, we are interested in the couplings of the fermion zero-modes with the first KK excitations of the gauge bosons. The wave functions for the nth KK gauge boson is given by [5, 6]:

$$\chi_n(y) = \frac{e^{ky}}{N_n} \left[J_1\left(\frac{m_n}{k}\, e^{ky}\right) + \alpha_n Y_1\left(\frac{m_n}{k}\, e^{ky}\right)\right]\,, \tag{2.7}$$

where m_n is the mass of the nth KK excitation of the gauge boson, N_n is the normalization, J_1 and Y_1 are Bessel functions, and the constant α_n is defined by $\alpha_n = -J_0(\frac{m_n}{k})/Y_0(\frac{m_n}{k}) = -J_0(\frac{m_n\, e^{k\pi R}}{k})/Y_0(\frac{m_n\, e^{k\pi R}}{k})$, which also determines the KK masses m_n.

For the zero mode fermions, one obtains the wave-functions

$$f_0^{L,R}(y) = \sqrt{\frac{k\pi R(1 \mp 2c_{L,R})}{e^{(1\mp 2c_{L,R})k\pi R} - 1}}\, e^{\mp c_{L,R}\, ky}\ . \tag{2.8}$$

However, it is more useful to consider the y-dependence of the kinetic terms of the KK fermions as the effective fermion wave-functions. In this case the y dependence is

$$\hat{f}_{L,R}(y) = \frac{e^{(\frac{1}{2} \mp c_{L,R})ky}}{N_{L,R}}\,, \tag{2.9}$$

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where we defined the normalization factors

$$\frac{1}{N_{L,R}} \equiv \sqrt{\frac{1 \mp 2c_{L,R}}{e^{k\pi R(1\mp 2c_{L,R})} - 1}} \tag{2.10}$$

Then, left-handed (right-handed) fermions with $c_L > 1/2$ ($c_R < -1/2$) are localized towards the Planck brane, whereas left-handed (right-handed) fermions with $c_L < 1/2$ ($c_R > -1/2$) are localized close to the TeV brane. Light fermions are of the first kind, while heavier fermions must be localized near the TeV brane. This is the case with the top quark, and now in this model also with all the fourth-generation fermions.

For values of the bulk mass parameter $c_L > 1/2$, the zero-mode fermion couples universally, as well as weakly, to the first KK gauge boson. For $c_L < 1/2$, the coupling can be considerably enhanced above the gauge coupling. We will consider that the fourth generation has bulk mass parameters that localize it very close to the TeV brane, and therefore it will have very strong couplings to the KK gauge bosons.

In addition to the strong interactions among fourth-generation zero-modes induced by the KK gauge bosons, there are interactions induced by bulk higher-dimensional operators. Of particular interest are the four-fermion bulk operators

$$\int dy\,\sqrt{g}\,\frac{C^{ijk\ell}}{M_P^3}\,\bar{\Psi}^i_L(x,y)\Psi^j_R(x,y)\bar{\Psi}^k_R(x,y)\Psi^\ell_L(x,y)\,, \tag{2.11}$$

where $C^{ijk\ell}$ are generic coefficients, with i, j, k, ℓ standing for generation indices as well as other indices such as isospin, and the $\Psi(x,y)$'s can be bulk quarks or leptons. The naive dimensional analysis (NDA) estimate of the 5D coefficients in (2.11) gives

$$C^{ijk\ell} \sim \frac{36\,\pi^3}{N}\,, \tag{2.12}$$

with N the number of fermion flavors that can be accommodated in a loop. If we assume that all the C's are of the same order, then $N \sim O(100)$.[1] These bulk operators lead to 4D four-fermion operators involving the various fermion zero and KK modes. The four-fermion interactions induced among fermion zero-modes are given by [5]

$$C^{ijk\ell}\,\frac{k}{M_P^3}\,\frac{e^{k\pi R(4-c^i_L-\tilde{c}^j_R-\tilde{c}^k_R-c^\ell_L)}-1}{4-c^i_L-\tilde{c}^j_R-\tilde{c}^k_R-c^\ell_L}\,\frac{\bar{\psi}^{i(0)}_L\psi^{j(0)}_R\bar{\psi}^{k(0)}_R\psi^{\ell(0)}_L}{N^i_L\,N^j_R\,N^k_R\,N^\ell_L}\,, \tag{2.13}$$

where we defined $\tilde{c}^i_R = -c^i_R$ for convenience.

Once again, localization is determining the size of these contributions. For instance, for bulk mass parameters for two of the fermions satisfying $(c^i_L, \tilde{c}^j_R) > 1/2$, i.e. two of the fermions being Planck localized, these four-fermion operators are exponentially suppressed. On the other hand if all four-fermions have bulk mass parameters localizing the zero-modes towards the TeV brane ($(c_L, \tilde{c}_R) < 1/2$), the corresponding contribution will be only suppressed by the TeV scale. In particular, the four-fermion interactions induced among

[1] For instance, if all the $C's$ are exactly equal, the number of fermions inside a loop mediating any given four-fermion interaction is $N = 80$.

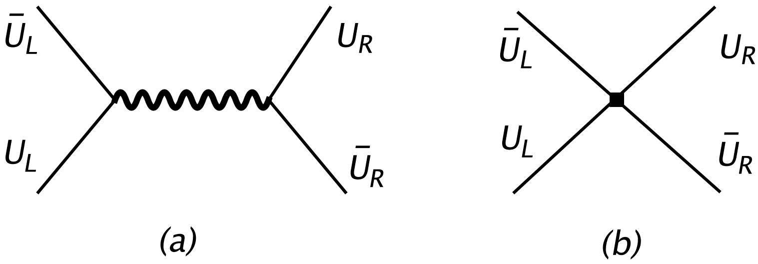

Figure 1: Two contributions to four-fermion interactions of the up-type fourth-generation quark: (a) from the interactions with a KK gluon; (b) from the four-fermion interactions induced by the bulk operators of (2.11).

fourth-generation zero-modes are only TeV suppressed, with a dimensionless coupling of the form

$$\sim C^{4444}\left(\frac{k}{M_P}\right)^3 \frac{(1-2c_L^4)(1-2\tilde{c}_R^4)}{2\,(2-\tilde{c}_R^4-c_L^4)}\,. \tag{2.14}$$

were the sub-indices in the coefficients denote flavor quantum numbers.

2.2 Four-fermion interactions and electroweak symmetry breaking

We now examine the four-fermion interactions among fermions induced by the exchange of KK gauge bosons. The strongest coupling of the fourth generation is that of the first KK gluon to the fourth generation quarks. For instance, considering the zero-mode U quark we have the following four-fermion interaction below the mass of the first excitation of the KK gluon, M_{KK}:

$$-\frac{g_{01}^L\,g_{01}^R}{M_{KK}^2}\,\left(\bar{U}_L\gamma_\mu T^A U_L\right)\,\left(\bar{U}_R\gamma_\mu T^A U_R\right)\,, \tag{2.15}$$

where U is the zero mode of the fourth-generation up-type quark, g_{01}^L and g_{01}^R are the left-handed and right-handed U couplings to the first KK gluon excitation, and T^A are the usual $\mathrm{SU}(3)_c$ generators. After Fierz rearrangement, we can re-write this interaction as

$$\frac{g_{01}^L\,g_{01}^R}{M_{KK}^2}\,\left\{\bar{U}_L^a U_R^a\;\bar{U}_R^b U_L^b-\frac{1}{N_c}\,\bar{U}_L^a U_R^b\;\bar{U}_R^b U_L^a\right\}\,, \tag{2.16}$$

where a, b are $\mathrm{SU}(3)_c$ indices. The color singlet term in (2.16) is attractive, whereas the color octet is repulsive, as well as suppressed by $1/N_c$. We then concentrate in the color singlet four-fermion interaction. Likewise, electroweak KK gauge bosons give similar, although much smaller, contributions. We will neglect them in what follows. However, we must take into account the contributions generated from bulk higher-dimensional operators such as (2.11), since they are generally comparable to the ones obtained from KK exchange. Both these contributions, depicted in figure 2.1, result in an effective four-fermion interaction of the fourth-generation quarks. For the U quark, for instance, we then have an effective four-fermion coupling

$$g_U^2 \equiv g_{01}^L\,g_{01}^R + x_1^2\,C_{uu}^{4444}\left(\frac{k}{M_P}\right)^3\frac{(1-2c_L^4)(1-2\tilde{c}_R^4)}{2(2-c_L^4-\tilde{c}_R^4)}\,, \tag{2.17}$$

where $x_1 \equiv M_{KK}/\Lambda_{\rm TeV}$, is the mass of the first KK gauge bosons in units of the TeV scale defined as $\Lambda_{\rm TeV} = k\, e^{-k\pi R} \sim O(1)$ TeV.

There is a value of g_U^2 above which a condensate forms

$$\langle \bar{U}_L\, U_R \rangle \neq 0\,, \tag{2.18}$$

leading to electroweak symmetry breaking and dynamical masses for the condensing fermions. It is possible to obtain the criticality condition on the coupling from a gap equation. Here we want to describe electroweak symmetry breaking through the vacuum expectation value of the ensuing composite scalar, the Higgs. For this purpose, we start from the Lagrangian

$$\mathcal{L} = \bar{U}_L i\, \not{D} U_L + \bar{U}_R i\, \not{D} U_R + \frac{g_U^2}{M_{KK}^2}\, \left(\bar{U}_L U_R\, \bar{U}_R U_L\right)\,, \tag{2.19}$$

This can be re-written as

$$\mathcal{L} = \bar{U}_L i\, \not{D} U_L + \bar{U}_R i\, \not{D} U_R + g_U\, \bar{Q}_L H U_R - M_{KK}^2 H^\dagger H + \text{h.c.}\,, \tag{2.20}$$

where $Q_L^T \equiv (U_L\; D_L)^T$, H is a non-propagating SU$(2)_L$ doublet, and we have omitted the down-type quark kinetic terms. At scales $\mu < M_{KK}$, H develops a kinetic term as well as a self-coupling, resulting in

$$\begin{aligned}\mathcal{L}(\mu) &= Z_{U_L} \bar{U}_L i\, \not{D} U_L + Z_{U_R} \bar{U}_R i\, \not{D} U_R + \cdots + Z_{g_U} g_U\, \bar{Q}_L H U_R + h.c. \\ &\quad + Z_H (D_\mu H)^\dagger D^\mu H - m_H^2 H^\dagger H - \frac{\lambda}{2}\left(H^\dagger H\right)^2\,,\end{aligned} \tag{2.21}$$

where the wave-function renormalizations Z_{U_L}, Z_{U_R}, Z_H,..., as well as Z_{g_U}, m_H and λ, can be easily computed in the one loop approximation. For instance the dominant contribution to m_H results in

$$m_H^2 = M_{KK}^2 \left(1 - \frac{g_U^2\, N_c}{8\pi^2}\right) + \cdots\;. \tag{2.22}$$

Thus, we see that the effective potential for H at low energies develops a non-trivial vacuum if

$$g_U^2 > \frac{8\pi^2}{N_c}\;. \tag{2.23}$$

This condition is easily satisfied in these models, even in the absence of the four-fermion operators of eq. (2.11), by giving enough localization to the fourth-generation. For instance, if $(c_L^4, \tilde{c}_R^4) < 0$ the KK-gluon induced four-fermion interactions are always super-critical. In addition, if we include the effects of KK fermions in the effective Higgs theory, the resulting critical coupling would be lower than the one obtained in (2.23). In any case, the exact value of the critical coupling is not important for the calculation of the spectrum in this model.

Equation (2.23) coincides with the criticality condition obtained by making a one-loop gap equation analysis of (2.19). Then, we see that if the couplings of zero-mode fermions to KK gauge bosons are strong enough, they could lead to electroweak symmetry breaking. Among the SM fermions, the best candidate for accomplishing this is the top quark, as

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in top-condensation models [1, 15, 16]. However, even if we assumed that the effective four-fermion interactions of top quarks were super-critical, this would lead either to a top mass that is too large, or to a cutoff that has to be above 10^{15} GeV. Earlier attempts to embed top-condensation in flat [17] and warped [18] extra-dimensional theories, required the condensation of a large number of KK fermions in order to obtain the correct value of m_t. In the present AdS$_5$ setup this is very difficult to achieve and requires unnaturally large values for the coefficients in (2.11), given that higher KK modes have weaker couplings. A fourth generation with zero modes highly localized towards the TeV brane is guaranteed to condense. For simplicity, we will consider here the case where only the up quark U condenses. The case with the D also condensing leads to a more complicated scalar sector [19].

The coefficient of the kinetic term of H, and its self-coupling, computed at one loop, are given by

$$Z_H = \frac{g_U^2 N_c}{16\pi^2} \ln\left(\frac{M_{KK}^2}{\mu^2}\right) , \tag{2.24}$$

$$\lambda = \frac{g_U^4 N_c}{8\pi^2} \ln\left(\frac{M_{KK}^2}{\mu^2}\right) , \tag{2.25}$$

where we have only included the up quark zero mode contributions. We notice that both Z_H and λ vanish at the cutoff $\Lambda = M_{KK}$, reflecting the compositeness conditions. Completing the renormalization procedure, we consider the scalar contributions to Z_{U_L} and Z_{U_R}, as well as the coupling renormalization Z_{g_U} coming from scalar exchange. After the replacements

$$Z_{U_L}^{1/2} U_L \to U_L\,, \qquad Z_{U_R}^{1/2} U_R \to U_R\,, \qquad Z_H^{1/2} H \to H\,, \tag{2.26}$$

and the definition of the renormalized quantities

$$\bar{m}_H^2 = \frac{m_H^2}{Z_H}, \qquad \bar{\lambda} = \frac{\lambda}{Z_H^2}, \qquad \bar{g}_U = \frac{Z_{g_U}}{\sqrt{Z_{U_L}\, Z_{U_R}\, Z_H}}\, g_U\,, \tag{2.27}$$

the renormalized lagrangian reads

$$\begin{aligned}\mathcal{L}_r &= \bar{U}_L i \not{D} U_L + \bar{U}_R i \not{D} U_R + \cdots + \bar{g}_U\, \bar{Q}_L H U_R + h.c. \\ &\quad + (D_\mu H)^\dagger D^\mu H - \bar{m}_H^2 H^\dagger H - \frac{\bar{\lambda}}{2}\left(H^\dagger H\right)^2 .\end{aligned} \tag{2.28}$$

Assuming that the criticality condition (2.23) is satisfied, the Higgs field H acquires a VEV

$$\langle H \rangle = \begin{pmatrix} v/\sqrt{2} \\ 0 \end{pmatrix} , \tag{2.29}$$

giving the condensing fermion a dynamical mass $m_U = \bar{g}_U v/\sqrt{2}$. Here we take $v \simeq 246$ GeV, which results in the correct value for M_W at this order in perturbation theory. The Higgs mass, $m_h = \sqrt{\bar{\lambda}} v$, can be computed in this approximation and satisfies the Nambu-Jona-Lasinio (NJL) relation, $m_h = 2\, m_U$. However, this simplistic prediction receives important corrections that will be addressed in the next section. The same can be said of the prediction for the dynamical fermion mass m_U, which at this level of accuracy must satisfy the Pagels-Stokar expression

$$v^2 = m_U^2 \frac{N_c}{8\pi^2} \ln\left(\frac{M_{KK}^2}{m_U^2}\right) , \tag{2.30}$$

which points to dynamical masses in the few hundred GeV for the up-type fourth generation quark. We will refine the predictions for the dynamical fermion masses and the Higgs mass in section 3, where we will make use of the full renormalization group evolution of the couplings $\bar{g}_U$ and $\bar{\lambda}$. But before that, we will address the origin of the masses of all other (non-condensing) fermions.

2.3 Fermion masses

In the previous section we have shown that if the zero modes of quarks of the fourth generation are localized enough, they can condense and break the electroweak symmetry. The condensate $\langle \bar{U}_L U_R \rangle$ results in a dynamical mass for the U quark zero mode. Higher-dimensional operators suppressed by M_P such as (2.11), will result in masses for all other zero-mode fermions upon condensation of the fourth-generation up-type zero-mode quark U, as well as of any other condensing fermion. For instance, we could imagine that flavor was a gauge symmetry and it was broken at the Planck scale. The same might be true of other symmetries, which may couple quarks and leptons at the very high scale.

When two of the zero-mode fermions are condensing, for instance U_L, and U_R in the model we are considering here, the operator in (2.11) results in masses for the other zero modes. This corresponds to $k = \ell = 4$. Light fermion masses result from fermions with bulk mass parameters $(c_L^i, \tilde{c}_R^j) > 1/2$. Assuming the condensate satisfies $\langle \bar{U}_R U_L \rangle \sim m_U^3$, these are

$$\begin{aligned} m_{ij} = C^{ij44} \left(\frac{k}{M_P}\right)^3 \left(\frac{m_U}{\Lambda_{\rm TeV}}\right)^2 & \frac{\sqrt{(2c_L^i - 1)(2\tilde{c}_R^j - 1)}\sqrt{(1-2c_L^4)(1-2\tilde{c}_R^4)}}{4 - c_L^i - \tilde{c}_R^j - \tilde{c}_R^4 - c_L^4} \\ & \times e^{k\pi R(1-c_L^i-\tilde{c}_R^j)}\, m_U \ . \end{aligned} \tag{2.31}$$

The masses in (2.31) are exponentially suppressed and lead to light fermion masses. On the other hand, for $(c_L^i, \tilde{c}_R^j) < 1/2$, we arrive at

$$m_{ij} = C^{ij44} \left(\frac{k}{M_P}\right)^3 \left(\frac{m_U}{\Lambda_{\rm TeV}}\right)^2 \frac{\sqrt{(1-2c_L^i)(1-2\tilde{c}_R^j)}\sqrt{(1-2c_L^4)(1-2\tilde{c}_R^4)}}{4 - c_L^i - \tilde{c}_R^j - \tilde{c}_R^4 - c_L^4}\, m_U\, , \tag{2.32}$$

which is un-suppressed and of order m_U, up to factors of $O(1)$. This is the case for the top quark, the fourth-generation quark D, as well as the fourth-generation leptons. Thus, all fourth-generation fermions have masses in the several hundred GeV, with their exact values depending on the details of their localization near the TeV brane. This picture of fermion masses is consistent with the one obtained with a TeV-brane-localized Higgs [5, 13]. Then, in the effective theory with a composite Higgs field described in the previous section, it is possible to obtain the observed 4D Yukawa couplings starting from these four-fermion interactions.

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A more precise prediction for the fourth-generation masses, as well as for the mass of the Higgs boson, can be obtained by considering the full renormalization group running. We do this in the next section.

3. Renormalization group effects and mass predictions

In order to obtain better predictions for the spectrum of the theory at low energies, we must consider the effects of the renormalization group running, especially on the Yukawa coupling of the fourth-generation up quark U, and the (renormalized) Higgs self-coupling $\bar{\lambda}$. Here we follow closely Bardeen, Hill and Lindner [1, 20].

3.1 Yukawa running and dynamical fermion mass

Considering the one-loop contributions of the Yukawa coupling $\bar{g}_U$ to the wave-function renormalizations Z_{U_L} and Z_{U_R}, as well as the contributions to Z_{g_U} and Z_H, one obtains, neglecting gauge interactions,

$$\frac{d\bar{g}_U}{dt} = \frac{\bar{g}_U^3}{16\pi^2}\left[\frac{3}{2} + N_c\right] , \tag{3.1}$$

where $t = \ln(\mu)$ and μ is the renormalization scale. At high energies, the Yukawa coupling blows up. At low energies, however, the gauge contributions are important. If we take them into account we have

$$\frac{d\bar{g}_U}{dt} = \frac{1}{16\pi^2}\left[\frac{9}{2}\,\bar{g}_U^3 - C(t)\,\bar{g}_U\right] , \tag{3.2}$$

where the $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ running couplings are taken into account in

$$C(t) = 8g_s^2(t) + \frac{9}{4}g^2(t) + \frac{17}{12}g'^2(t) , \tag{3.3}$$

and we will use the values of the couplings as extracted in the particle data book [21]. Equation (3.2) is solved with the boundary condition

$$\bar{g}_U \to \infty, \qquad \text{for} \quad \mu \to \Lambda , \tag{3.4}$$

where Λ is the cutoff, and we take $\Lambda = M_{KK}$. The main effect is from the QCD coupling $g_s(t)$. The solution for the physical mass m_U comes from $m_U = \bar{g}_U(m_U)\,v/\sqrt{2}$. In figure 2 we show the result for the dynamical fourth-generation mass m_U as a function of the cutoff Λ. Since the cutoff is $\Lambda = M_{KK} = O(\text{few})$ TeV, we predict the dynamical fourth-generation mass in the range $m_U \sim (600 - 700)$ GeV, somewhat lower than the naive prediction of the previous section in (2.30). Potentially large mixing of fourth-generation zero-mode fermions with their KK modes might lower their masses even further, perhaps as much as 30% [22].

The masses of the other fourth-generation fermions depend on the choices of their bulk mass parameters, but they are typically of the order of m_U. Thus, we will be able to choose the amount and sign of isospin violation introduced by the zero-mode fourth generation.

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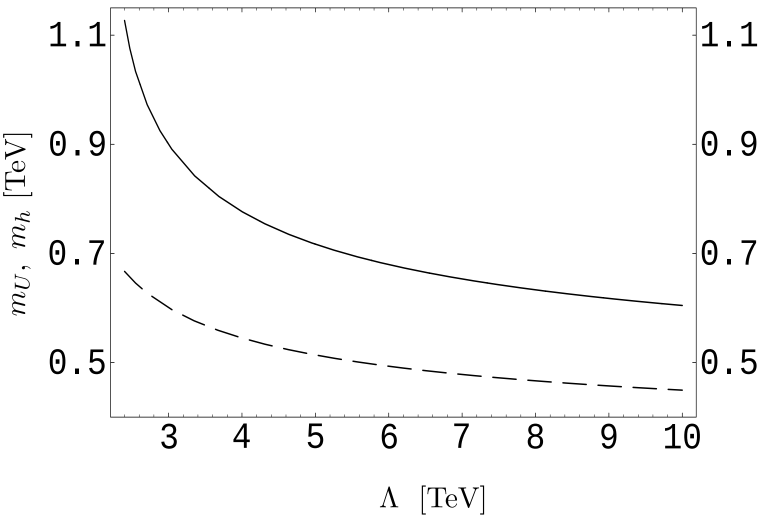

Figure 2: Dynamical mass of the fourth-generation up quark, m_U (dashed-line); and the physical Higgs mass m_h (solid line), both vs. the cutoff Λ, in TeV.

3.2 The Higgs mass

The renormalized Higgs self-coupling determines the Higgs mass through $m_h = \sqrt{\bar{\lambda}(m_h)}\, v$. If we neglect the gauge interactions and only consider the effect of the fourth and third generation Yukawa couplings, the renormalization group evolution of $\bar{\lambda}$ is given by

$$\frac{d\bar{\lambda}}{dt} = \frac{12}{16\pi^2}\left[\bar{\lambda}^2 + \sum_f \bar{g}_f^2\bar{\lambda} - \sum_f \bar{g}_f^4\right], \tag{3.5}$$

where $\bar{g}_f$ is the Yukawa coupling of the fermion f, the sums are over the fourth-generation quarks and a lepton doublet, and also include the top quark. From the expressions (2.25) and (2.27) for $\bar{\lambda}$, one would only obtain the last two terms in (3.5). However, from these expressions we also see that $\bar{\lambda}$ diverges at the cutoff $\Lambda = M_{KK}$. It is then consistent to consider the presence of the tree-level $\bar{\lambda}$ interaction that gives rise to the first term in (3.5). This is then the same RGE as for the SM Higgs self-coupling.

The solutions of (3.5) must satisfy compositeness conditions, as determined by eqs. (2.24) and (2.25). We may cast these by making the replacement $H \to \bar{g}_U\, H$ in (2.28), the renormalized lagrangian. This results in an effective Higgs self-coupling that goes like $\bar{\lambda}/\bar{g}_U^4$, which should go to zero at the cutoff Λ, to satisfy compositeness. Then, we see that $\bar{\lambda}$ must diverge slower than $\bar{g}_U^2$. This implies that the solutions to (3.5) flow to an ultra-violet fixed-point [1], such that

$$\bar{\lambda} \simeq \bar{g}_U^2\, x_+\,, \tag{3.6}$$

with $x_+ = (\sqrt{65}-1)/8 \sim 0.88$, and where we have considered that all the fourth-generation Yukawas are of order $\bar{g}_U$. Making use of our results for $\bar{g}_U$, this results in

$$m_h \simeq 1 \text{ TeV}\,, \tag{3.7}$$

for a cutoff $\Lambda \sim 2.5$ TeV, which is still very close to the naive NJL prediction.

Considering the electroweak gauge corrections, the full RGE for $\bar{\lambda}$ now is

$$\frac{d\bar{\lambda}}{dt} = \frac{12}{16\pi^2}\left[\bar{\lambda}^2 + \left(\sum_f \bar{g}_f^2 - A(t)\right)\bar{\lambda} + B(t) - \sum_f \bar{g}_f^4\right] , \tag{3.8}$$

where

$$\begin{aligned} A(t) &= \frac{1}{4}g'^2(t) + \frac{3}{4}g^2(t) \\ B(t) &= \frac{1}{16}g'^4(t) + \frac{1}{8}(g(t)g'(t))^2 + \frac{3}{16}g^4(t) . \end{aligned} \tag{3.9}$$

As we can see from the figure 2, the addition of the gauge contributions does not modify the prediction for m_h greatly. This remains a very heavy Higgs, if the cutoff is kept not far above the TeV scale.

4. Electroweak precision constraints

Bulk RS models on which we based our construction, have an enlarged isospin symmetry given by the extension from the SM gauge group to $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_X$. This forbids tree-level contributions to the T parameter. On the other hand, there are important contributions to the S parameter already at tree-level [13]. These can be seen as coming from the interactions of light (Planck-localized) fermions with the gauge bosons, through the KK modes. The modified couplings, being universal, can be re-absorbed into a redefinition of the gauge fields, resulting in contributions to the oblique parameters S and T. Particularly dangerous is the S parameter contribution, given by

$$S_{\text{tree}} \simeq 12\pi \,\frac{v^2}{M_{KK}^2} . \tag{4.1}$$

Additional tree-level contributions correspond to operators of dimension eight or higher, and are further suppressed by factors of v^2/M_{KK}^2.

In the present model, the presence of a fourth generation induces additional contributions to electroweak observables at one loop, both from the fourth-generation zero modes, as well as their KK excitations. The presence of a degenerate SM fourth generation (the zero-modes) results in a positive shift of the S parameter given by

$$S_{4g} \simeq \frac{2}{3\pi} , \tag{4.2}$$

with this result somewhat smaller if the down sector is lighter than the up. Recent re-examination of the constraints on a fourth generation coming from electroweak precision measurements has shown [23] that the presence of these states is not in serious contradiction with data, as it is concluded in ref. [21]. This is particularly the case if the fourth-generation quarks have splittings giving a positive contribution to T. In our model, this can be naturally achieved by having the up-type quark more localized than the down-type such that $m_U > m_D$. In realizations where only the up quark condenses this is most easily achieved, but it can be also the case even if both the U and the D condense.

The contributions of the KK fermions to S and T can be summed up. Their calculation is cumbersome and we will leave it for a future publication, where we will put together all the electroweak precision constraints of the model [22]. However, we can already conclude that their contribution will be under control and smaller than the one from the zero-modes.

Also, the fact that the Higgs is heavy results in a positive shift of the S parameter. The standard one loop contribution to S from a heavy Higgs, results in

$$\Delta S^h_{SM} \simeq \frac{1}{12\pi} \ln \left(\frac{m_h}{m_h^{\text{ref.}}} \right)^2 , \tag{4.3}$$

Taking the reference value to be $m_h^{\text{ref.}} = 114$ GeV, results in a $\Delta S \simeq +0.1$ for the typical Higgs mass in our model, which corresponds to a shift of the origin of the $S-T$ plot.[2] The 95% C.L. bound [21] for a heavy Higgs is $S < 0.09$. Thus, it is difficult to accommodate the positive S contributions (4.1) and (4.2) with a heavy Higgs. For instance, for $M_{KK} \simeq 3$ TeV we have $S_{\text{tree}} \simeq 0.25$, and typically $S_{4\text{g}} \lesssim 0.2$. Then, we conclude that in its present form the model appears to be disfavored by current electroweak precision constraints.

However, given that the theory is strongly coupled, care must be taken before drawing definite conclusions from leading order perturbative calculations. For instance, it is possible to show that the strong coupling of the Higgs to the KK vector resonances results in a reduction of the Higgs contribution to S. A similar effect takes place with the contribution of the fourth generation if, as is the case here, it is strongly coupled to the KK vectors. A thorough study of these strong coupling effects will be performed elsewhere [22].

5. Phenomenology

The class of models presented here has a very rich phenomenology at the LHC. Some of its aspects will depend on the specific realization of the fourth-generation condensation model. For instance, the scalar sector could be richer if the fourth-generation down quark condenses, leading to a two-Higgs doublet spectrum. Also, the choice of fermion assignment to the $\mathrm{SU}(2)_R$ group results in at least two possibilities for the spectrum of relatively light KK fermions. However, there are some generic features that would constitute signals for these class of models. If the zero-mode spectrum constitutes a complete fourth generation, its discovery at the LHC, in association with a heavy Higgs, would give a hint that the fourth-generation could be associated with electroweak symmetry breaking. More definite proof of this, would be the observation of the strong coupling of the fourth-generation quarks to the KK gluon excitations. In what follows we briefly discuss some generic phenomenological features of the model discussed in the previous sections.

The production cross section of fourth-generation quark pairs is of about [25]

$$\begin{aligned} &\sigma_{Q_4\bar{Q}_4} \simeq 1\ pb, &&\text{for} &&m_{Q_4} = 600 \text{ GeV}\,, \\ &\sigma_{Q_4\bar{Q}_4} \simeq 0.1\ pb, &&\text{for} &&m_{Q_4} = 900 \text{ GeV}\,. \end{aligned}$$

[2] An additional cutoff-dependent contribution to S appears in models where the Higgs is a pseudo-Nambu-Goldstone boson, as pointed out in ref. [24].

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Thus, approximately 1000 events per quark type will be produced in a typical low luminosity year for a 900 GeV fourth-generation quark. However, the reach could be limited to masses below this due to backgrounds.

If $m_U > m_D$, as we considered here in order to have only the U quark condense, then for $(m_U - m_D) > M_W$, the up-type quark would decay as $U \to DW$. The down-type quark would decay almost exclusively to the top quark through $D \to tW$. Thus, the pair production of U pairs results in the decay chain $U\bar{U} \to W^+W^-W^+W^-W^+W^-b\bar{b}$, with six W's plus two b jets. This signal has not been studied at the LHC and it appears challenging due to the large number of jets. However, it appears that it might be possible to device a way to reconstruct the D quarks, since we could use the leptonic decay of a W from a U decay for triggering. On the other hand, the D pair production results in the chain $D\bar{D} \to W^+W^-W^+W^-b\bar{b}$, which has been studied in ref. [25]. If $(m_U - m_D) < M_W$, then U would decay through $U \to bW$. Then, $U\bar{U}$ production is identical to top pair production with the exception of the quark mass. A preliminary study in ref. [25] shows that with $100\, fb^{-1}$ luminosity it is possible to have a significant signal above background for masses up to at least 700 GeV. Other decay modes, involving significant mixing with the third generation quarks, are studied in ref. [26].

Regarding leptons, the standard production cross section for a pair of charged leptons $L\bar{L}$ or of massive neutrinos $N\bar{N}$, is much smaller than in the quark case, since these are electroweak processes. Typically, for $m_L, m_N \simeq 700\,\text{GeV}$, cross sections are a few fb. For instance, if $(m_L - m_N) > M_W$, the charged lepton could decay through $L \to N_L W$. The left-handed neutrino would then decay trough mixing with the lighter generations, through $N_L \to \ell W$, with $\ell = \tau, \mu, e$. If these mixings are small enough, the decay might occur outside the detector, leading to a large missing E_T signal. If, on the other hand, $(m_L - m_N) < M_W$, then the charged lepton also must decay through mixing with lighter leptons, as in $L \to \nu W$. Once again, if the intergenerational mixings of the fourth-generation leptons are small, this could result in a slow charged track in the detector, which can be easily identified and might even allow the measurement of the charged lepton mass [25]. Finally, if we assume the existence of a right-handed neutrino zero-mode N_R, for instance in order to obtain a Dirac mass for the fourth-generation neutrino zero-mode, then its production and decay will depend on the transformation properties of the bulk field it belongs to. For instance, if this transforms as a $(\mathbf{1}, \mathbf{3})_\mathbf{0}$, it only couples to the KK excitation of the Z', the state orthogonal to the SM Z. Thus, its production cross section is rather small. Its decays are also suppressed since they can only proceed through a three body decay further suppressed by the probably small mixings with the lighter lepton generations. These events could have a very characteristic signal of large missing E_T and little activity in the central region, albeit very rare. On the other hand, if the bulk field resulting in a fourth-generation right-handed neutrino transforms as a $(\mathbf{1}, \mathbf{1})_\mathbf{0}$, the only couplings of the zero-mode are through the four-fermion operators of (2.11) and the effective Yukawa coupling they generate. Then, the operators responsible for their production and decay are effectively Planck-suppressed, making them possibly stable in cosmological time scales.

But the most distinct signal of the model will not be the presence of a heavy fourth generation in combination with a heavy Higgs. In order to clearly detect this class of

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models, one must prove that the fourth-generation is strongly coupled to the TeV scale resonances, the gauge KK modes responsible for the condensation of the fourth-generation quarks. The main signal for this is the production of fourth-generation quarks and leptons through s-channel production of the KK gauge bosons. In particular, the produced KK gauge bosons, if strongly coupled to the fourth-generation, would decay to it preferentially. Then, for the KK gluon for instance, we have that

$$\frac{\mathrm{Br}(G^{(1)} \to U\bar{U})}{\mathrm{Br}(G^{(1)} \to t\bar{t})} \sim (5-10)\,, \tag{5.1}$$

depending of the parameters of the model. A careful study of the possibility of reconstructing these very high-invariant mass events must be done in order to evaluate how well can this signal be seen at the LHC. On the other hand, the contact four-fermion interactions coming from (2.11), such as $q\bar{q}U\bar{U}$, are much more suppressed, typically by the light quark masses.

Finally, we note that the spectrum of KK fermions includes states that typically have masses not very different from the fourth-generation zero-modes'. Some of these should have very different signals compared to a standard fourth-generation, given their exotic quantum numbers [27]. In general, a very detailed study of all these signals, and the corresponding backgrounds, must be carried out in order to assess the reach of the LHC in this model. We leave this for future work [22].

6. Conclusions and outlook

We have shown a mechanism for the breaking of the electroweak symmetry and the generation of fermion masses through the condensation of a fourth generation. In the context of a 5D theory in a slice of AdS_5, the super-critical interactions of the fourth-generation zero-mode quarks are induced by the KK excitations of the gluon, as well as by bulk higher-dimensional operators. These are strong due to the localization of the fourth-generation zero-modes close to the IR brane. The condensation of the fourth-generation quarks leads to electroweak symmetry breaking and results in a heavy Higgs, with a mass $m_h \simeq 900$ GeV, for a KK mass of about 2.5 TeV. The unitarization of SM amplitudes is achieved partially by this heavy Higgs, and partially by the presence of the KK gauge bosons. The condensing quarks, the zero-mode of the fourth-generation up-quark sector, acquires a dynamical mass of about $m_U \simeq (600-700)$GeV for the same value of the KK mass. Larger values of the KK gauge masses result in lighter Higgs and dynamical fermion masses. However, as the KK mass is increased the theory becomes more fine-tuned.

Fermion masses for the lighter three generations, as well as the non-condensing fourth-generation fermions, are generated by higher-dimensional bulk operators suppressed by the Planck mass. After dimensional reduction, these result in four-fermion interactions amongst zero-mode fermions. The ones involving two condensing quarks will give rise to mass terms for the remaining two fermions. This results in the necessary Yukawa textures and the observed fermion masses and mixings. Thus, in this model the mechanism of electroweak symmetry breaking requires flavor violation in the bulk, and is intimately related to the fermion masses. This is to be contrasted with the proposal of ref. [28], where

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fermions are de-localized as a way to evade electroweak constraints; as well as with the one of ref. [29], where there is a flavor symmetry in the 5D bulk.

Regarding electroweak constraints, and as is the case with all bulk Randall-Sundrum models where the fermion localization naturally explains the fermion mass hierarchy, this model contains a tree-level contribution to the S parameter. In addition, the presence of a heavy Higgs results in a positive shift of the S parameter at one loop; and the fourth-generation zero-modes induce one loop contributions to both S and T. A full study of the electroweak precision constraints on the model, including the contributions from KK modes as well as effects coming from strong coupling, is left for a separate publication [22]. However, we can already conclude that the loop contributions to S are not the decisive factor given the presence of a tree-level contribution. In other words, the situation of Randall-Sundrum bulk models is not made significantly worse by the presence of a bulk fourth-generation.

The phenomenology of the model at the LHC involves the discovery of a strongly coupled heavy fourth generation, the signal for a heavy Higgs, typically associated with enhanced longitudinal gauge boson scattering. To the usual fourth generation production and decay, this model adds the presence of high invariant mass production of the fourth generation through its strong coupling to the KK gauge bosons, particularly the gluon. These signals combined would constitute strong evidence that the condensation of the fourth-generation quarks is the origin of electroweak symmetry breaking [22]. Other possible phenomenological consequences of the model are, among others, the modification of the Higgs production cross section and decay widths [23], flavor physics observables and possible effects in neutrino physics and astrophysics. All of them deserve further study.

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