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Darkening the little Higgs

Travis A.W. Martin *, Alejandro de la Puente

TRIUMF, 4004 Wesbrook Mall, Vancouver, BC V6T 2A3, Canada

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ABSTRACT

We present a novel new method for incorporating dark matter into little Higgs models in a way that can be applied to many existing models without introducing T-parity, while simultaneously alleviating precision constraints arising from heavy gauge bosons. The low energy scalar potential of these dark little Higgs models is similar to, and can draw upon existing phenomenological studies of, inert doublet models. Furthermore, we apply this method to modify the littlest Higgs model to create the next to littlest Higgs model, and describe details of the dark matter candidate and its contribution to the relic density.

1. Introduction

Little Higgs (LH) models [1–3] are extensions of the Standard Model (SM) that stabilize the electroweak scale with a light Higgs boson and weakly coupled new physics. These models resolve the fine-tuning problem within the SM by embedding the Higgs boson within a non-linear sigma field, and by introducing new gauge and fermion states that result in a collective breaking of the scalar Higgs potential. This collective symmetry breaking ensures cancellation of the quadratic divergences that result from radiative corrections from gauge boson and top quark loops that plague the SM Higgs boson.

The challenges in constructing a modern little Higgs model include: generating a natural mass hierarchy between the heavy top partner(s) and heavy gauge bosons that fits within precision electroweak constraints; avoiding the generation of a dangerous singlet in the scalar potential [4]; and, in light of the mounting evidence for dark matter, the inclusion of a dark matter candidate. For example, the littlest Higgs model [5–7] and simplest little Higgs model [8] do not include a dark matter candidate, and are largely constrained by precision measurements [9–11]. While the bestest little Higgs (BLH) model [12] resolves these precision constraint issues by including a custodial $SU(2)$ symmetry and introducing a second non-linear sigma field that couples only to the gauge bosons, it does not include a dark matter candidate.

It has been noted that certain classes of little Higgs models may contain discrete symmetries that can be used to introduce a viable dark matter candidate. In particular, three such classes of models have been studied: theory space models [13], T-parity models [14] and skyrmion models [15,16]. In the latter, T-parity [14] requires new fermions, forces the gauge couplings to be equal, $g_1^{(\prime)} = g_2^{(\prime)}$, forces conservation of a T-charge for all interactions (therefore, the lightest T-odd state is stable), and results in an elimination of the triplet vacuum expectation value (vev). Theory space models [13] contain a Z_4 symmetry that can be used to interchange the non-linear sigma model fields amongst themselves. Within this class of models, the scalar identified with the SM-like Higgs boson breaks the Z_4 symmetry down to a Z_2 symmetry after electroweak symmetry breaking (EWSB), and the lightest particle charged under the Z_2 may become a viable dark matter candidate. Additionally, dark matter can arise in some little Higgs models from topological considerations [15,16]. In these models, skyrmions take the form of topological solitons.

In this Letter, we explore an alternative method of introducing dark matter to little Higgs models by incorporating a second non-linear sigma field, Δ. This expands upon the concept introduced in the bestest little Higgs model [12] and in another T-parity model [17], and provides a relatively simple means of implementing an inert doublet potential [18–20] – in effect, we prescribe a means of little Higgs-ing the inert doublet models. It should be noted that this is not the only implementation of an inert doublet potential in little Higgs models (see [21]). This presents a new class of little Higgs models, dark little Higgs (DLH) models, which follow the general structure:

* Corresponding author.
E-mail addresses: tmartin@triumf.ca (T.A.W. Martin), adelapue@triumf.ca (A. de la Puente).

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- duplicate global symmetry (G_Δ/H_Δ duplicates group structure of G_Σ/H_Σ) that breaks at scale $F > f$;
- G_Δ gauged in the same way as G_Σ;
- and, fermions transform only under G_Σ.

Since fermions do not transform under the second global symmetry, the complex doublet embedded in Δ does not develop a non-zero vev, and thus remains as a possible dark matter candidate. Additionally, by following this prescription, the heavy top partner masses are disconnected from the mass of the heavy gauge bosons, which relaxes electroweak precision constraints on the models without reintroducing fine-tuning constraints.

In this Letter, we describe the details of a simplistic version of this by modifying the littlest Higgs model into the next to littlest Higgs model, a DLH class model, and explore the relic abundance generated by the inert doublet.

2. The model

The littlest Higgs is based on a non-linear sigma field (Σ) that parametrizes an $SU(5)_\Sigma/SO(5)_\Sigma$ coset space. We introduce a second non-linear sigma field, Δ, parametrizing a separate coset space, $SU(5)_\Delta/SO(5)_\Delta$, but require that both the $SU(5)_\Sigma$ and $SU(5)_\Delta$ global symmetries contain the same gauged $[SU(2)\times U(1)]^2$ subgroup. Fermions transform only under the $SO(5)_\Sigma$ symmetry, and so the scalar doublet embedded in Δ does not acquire a radiatively generated negative mass squared. As with other little Higgs models, this description does not explain the physics origin of the non-linear sigma model, which is relevant only at or above the "compositeness" scale $\lambda \sim 4\pi f$.

The $SU(5)_\Sigma$ symmetry is broken to $SO(5)_\Sigma$ at a scale f, as in the littlest Higgs, while $SU(5)_\Delta$ is broken to $SO(5)_\Delta$ at a scale F ($> f$). The vacuum expectation values that generate this breaking are the same as in the littlest Higgs model, given by:

$$\Sigma_0 = \begin{pmatrix} 0 & 0 & \mathbb{1}_{2\times2} \\ 0 & 1 & 0 \\ \mathbb{1}_{2\times2} & 0 & 0 \end{pmatrix}, \qquad \Delta_0 = \begin{pmatrix} 0 & 0 & \mathbb{1}_{2\times2} \\ 0 & 1 & 0 \\ \mathbb{1}_{2\times2} & 0 & 0 \end{pmatrix}. \quad (1)$$

The non-linear sigma fields are then parameterized as:

$$\Sigma(x) = e^{2i\Pi_\Sigma/f}\Sigma_0, \qquad \Delta(x) = e^{2i\Pi_\Delta/F}\Delta_0 \quad (2)$$

where $\Pi_\Sigma = \sum_a \pi^a_\Sigma X^a$ and $\Pi_\Delta = \sum_a \pi^a_\Delta X^a$, summing over the 14 Goldstone bosons ($\pi^a_{\Sigma,\Delta}$) corresponding to the 14 generators (X^a) in each sector. In the littlest Higgs model, four fields corresponding to four of the broken generators are eaten to give mass to the heavy gauge bosons, and three are eaten to give mass to the SM gauge bosons, leaving seven observable scalar states. In our model, there are 14 broken generators for each of the Σ and Δ sectors (total of 28), and a total of seven are eaten to give mass to the gauge bosons, leaving 21 observable scalars.

Both $SU(5)$ symmetries are gauged by the same $[SU(2)\times U(1)]^2$ subgroups, with generators $Y_1 = \mathrm{diag}(-3,-3,2,2,2)/10$ and $Y_2 = \mathrm{diag}(-2,-2,-2,3,3)/10$ for the two $U(1)$ groups, and

$$Q^a_1 = \begin{pmatrix} \sigma^a/2 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \qquad Q^a_2 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -\sigma^{a*}/2 \end{pmatrix} \quad (3)$$

for the two $SU(2)$ groups. In this notation, σ^a are the Pauli matrices.

The new fields of the Δ non-linear sigma field are embedded in the Pion matrix as:

$$\Pi_\Sigma = \begin{pmatrix} 0 & h^\dagger/\sqrt{2} & \phi^\dagger \\ \xi/\sqrt{2} & 0 & h^*/\sqrt{2} \\ \phi & h^T/\sqrt{2} & 0 \end{pmatrix} + \left(Q^a_1 - Q^a_2\right)\eta^a$$
$$+ \sqrt{5}(Y_1 - Y_2)\sigma,$$
$$\Pi_\Delta = \begin{pmatrix} 0 & \xi^\dagger/\sqrt{2} & \chi^\dagger \\ \xi/\sqrt{2} & 0 & \xi^*/\sqrt{2} \\ \chi & \xi^T/\sqrt{2} & 0 \end{pmatrix} + \left(Q^a_1 - Q^a_2\right)\alpha^a$$
$$+ \sqrt{5}(Y_1 - Y_2)\beta, \quad (4)$$

where ξ and χ are the analogous fields to the h and ϕ from the Σ sector, and the real triplet (η^a, α^a) and singlet (σ, β) representations of the two non-linear sigma fields mix to form a combination that becomes the longitudinal components of the heavy gauge bosons ($\alpha^a_e = (f\eta^a + F\alpha^a)/\sqrt{f^2+F^2}$ and $\beta_e = (f\sigma + F\beta)/\sqrt{f^2+F^2}$), and an orthogonal combination that is physical.

These new fields couple to the gauge bosons in the normal way, via the kinetic term of the Lagrangian, such that,

$$\mathcal{L}_K = \frac{f^2}{8}\mathrm{Tr}\left[(D_\mu\Sigma)\left(D^\mu\Sigma\right)^\dagger\right] + \frac{F^2}{8}\mathrm{Tr}\left[(D_\mu\Delta)\left(D^\mu\Delta\right)^\dagger\right]. \quad (5)$$

The covariant derivative is given as

$$D_\mu\Sigma(\Delta) = \partial_\mu\Sigma(\Delta) - i\sum_j g_j W^a_j\left(Q^a_j\Sigma(\Delta) + \Sigma(\Delta)Q^{aT}_j\right)$$
$$- i\sum_j g'_j B_j\left(Y_j\Sigma(\Delta) + \Sigma(\Delta)Y_i\right), \quad (6)$$

where the sum is over $j = 1, 2$ for each of the two $SU(2)\times U(1)$. The heavy gauge boson masses pick up an extra contribution proportional to F^2, such that $M^2_{W_H} = \frac{1}{4}(g^2_1 + g^2_2)(f^2 + F^2)$ and $M^2_{B_H} = \frac{1}{20}(g'^2_1 + g'^2_2)(f^2 + F^2)$.

The Coleman–Weinberg (CW) derived couplings (λ's) for the h and ϕ in the scalar potential remain predominantly unchanged at leading order, as factors of F cancel out, leaving a dependence only on the scale Λ. Factors of F still contribute in the μ^2 term, which contains logarithmic divergences, through the masses of the heavy gauge bosons. The negative contribution from the heavy quark sector is still dominant in the μ^2 term in the potential, and induces spontaneous symmetry breaking.

We can examine the degree of fine tuning in the model as in [22] by examining the logarithmically divergent contributions to the μ^2 term in the scalar potential. Examining $\delta_T\mu^2$, $\delta_W\mu^2$, $\delta_B\mu^2$ and $\delta_\phi\mu^2$, we similarly find that $\delta_T\mu^2$ is responsible for the largest degree of fine tuning of the μ parameter. For a Higgs boson mass of 125 GeV, and scale parameters $f = 1$ TeV and $F = 5$ TeV, we find $\delta_W\mu^2/m^2_h < 11$, as compared with $\delta_T\mu^2/m^2_h < 180$. Thus it is clear that the degree of fine tuning in the model is controlled by the heavy quark sector, and larger values of $M_{W'}$ that result in a relaxation of electroweak (EW) precision constraints are viable without significantly increasing the degree of fine tuning.

Other EW precision constraints arise in the model as a result of the triplet vev, v'. The scalar potential for ϕ is unchanged from the littlest Higgs model, which provides the relation $v' < (v/4f)v$ [23]. Since the v' contributions to the EW precision observables are subdominant over those proportional to v^2/f^2 (or $M^2_W/M^2_{W'}$) [23] for most of the parameter space, the overall constraints on the scales f and F arising from EW precision observables will be improved over the original littlest Higgs model. In [6], it was argued that v' passes the constraints on Δg^Z_1 for values of $v' < 10\%v$, which is easily satisfied within the NLH model.

The masses of the ϕ and h fields in the Σ sector are similar to those found in the littlest Higgs. The χ triplet obtains a quadratically divergent mass from the one loop CW potential, while the ξ doublet only obtains a logarithmically divergent mass. The dominant terms in the masses of these states are given by: