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Vacuum Stability and Radiative Electroweak Symmetry Breaking in an SO(10) Dark Matter Model

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Vacuum stability in the Standard Model is problematic as the Higgs quartic self-coupling runs negative at a renormalization scale of about 10^{10} GeV. We consider a non-supersymmetric SO(10) grand unification model for which gauge coupling unification is made possible through an intermediate scale gauge group, $G_{\text{int}} = \text{SU}(3)_C \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L}$. G_{int} is broken by the vacuum expectation value of a **126** of SO(10) which not only provides for neutrino masses through the seesaw mechanism, but also preserves a discrete \mathbb{Z}_2 that can account for the stability of a dark matter candidate, here taken to be the Standard Model singlet component of a bosonic **16**. We show that in addition to these features, the model insures the positivity of the Higgs quartic coupling through its interactions to the dark matter multiplet and **126**. We also show that the Higgs mass-squared runs negative triggering electroweak symmetry breaking. Thus the vacuum stability is achieved along with radiative electroweak symmetry breaking and captures two more important elements of supersymmetric models without low energy supersymmetry. The conditions for perturbativity of quartic couplings and for radiative electroweak symmetry breaking lead to tight upper and lower limits on the dark matter mass, respectively, and this dark matter mass region (1.35–2 TeV) can be probed in future direct detection experiments.

Introduction.—With the discovery of the Higgs boson at both the ATLAS [1] and CMS [2] detectors, the Standard Model (SM) of particle physics appears to be very well established. However, as yet, there is no verified explanation for neutrino masses, and the nature of dark matter (DM) remains elusive. Both signal the need for beyond the SM physics. The experimental value of the Higgs mass, $m_h = 125.09 \pm 0.24$ GeV [3], also points to new physics at some higher energy scale. The Higgs quartic self-coupling, λ , runs toward negative values at high energy. The lower limit on m_h to ensure the positivity of λ out to the Planck scale is 129.4 ± 1.8 GeV [4] which appears to be violated. DM is often introduced in supersymmetric extensions of the SM with R -parity [5]. In supersymmetric models, the problem of vacuum stability associated with the Higgs quartic coupling is avoided as the tree level coupling is determined by a combination of the gauge couplings and is positive definite. In addition, supersymmetric models offer a mechanism for triggering electroweak symmetry breaking via radiative effects [6].

In the absence of supersymmetry, the instability in the Higgs potential occurs at a renormalization scale of about 10^{10} GeV for $m_h \simeq 125$ GeV. Therefore one might anticipate new physics playing a role at this intermediate scale or below [7, 8], such as the seesaw mechanism for generating neutrino masses which has long been associated with an intermediate scale [9]. This mechanism is very naturally realized in SO(10) grand unified theories (GUTs) [10, 11] where the right-handed neutrino is included as the SM singlet component of the **16** of SO(10) and incorporates a full generation of matter fields in a single

representation. SO(10) contains several subgroups, G_{int} which contain the SM gauge group as a subgroup and it is well known that the symmetry breaking scale, M_{int} , of G_{int} can be determined by requiring that the gauge couplings unify at a single scale $M_{\text{GUT}} > M_{\text{int}}$ [11–15].

Stable dark matter can also be incorporated in SO(10) models in a straightforward way [13–18]. As a rank-five group, SO(10) includes an additional U(1) symmetry, which is assumed to be broken at the intermediate scale. If the Higgs field that breaks this additional U(1) symmetry belongs to a **126** dimensional representation, then a discrete \mathbb{Z}_2 symmetry is preserved at low energies [19]. If we restrict our attention to relatively small representations ($\leq \mathbf{210}$), the **126** Higgs field leaving a \mathbb{Z}_2 symmetry is the only possibility for a discrete symmetry [14, 20]. For example, a scalar dark matter candidate will be stabilized by the \mathbb{Z}_2 symmetry if it is a member of either a **16** or **144** representation.

In this letter, we show that the two aforementioned attributes of supersymmetric extensions to the SM, namely, vacuum stability and radiative electroweak symmetry breaking, are also natural consequences of SO(10) models with an intermediate scale gauge group. For definiteness, we consider here a SM singlet dark matter candidate originating from a single **16** of SO(10) as in model SA₃₂₂₁ in Ref. [15] based on the intermediate gauge group $\text{SU}(3)_C \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L}$. In this model, the intermediate scale is found to be $M_{\text{int}} \simeq 10^9$ GeV and is small enough to allow the couplings of the **126** Higgs field to the SM Higgs to lift the Higgs quartic coupling through the threshold corrections before it turns

negative. The presence of the singlet scalar DM at low energies also deflects the running of the Higgs quartic coupling. Moreover, we show that the negative mass-squared needed for electroweak symmetry breaking runs positive due the coupling of the Higgs field with the DM singlet.

The requirement for the radiative electroweak symmetry breaking imposes a lower bound on the DM–Higgs coupling. This then leads to a lower limit on the DM mass if one assumes that the thermal relic abundance of the DM agrees with the observed DM density $\Omega_{\text{DM}} h^2 \simeq 0.12$ [21]. On the other hand, perturbativity of the couplings in the model gives an upper limit on the DM–Higgs coupling, and thus on the DM mass. As a result, a finite DM mass region is allowed by these two conditions. We find that this mass range can be probed in the XENON1T experiment [22].

An exemplary $SO(10)$ model with stable dark matter.—When one combines the number of possible intermediate scale gauge groups with the multitude of choices for dark matter and Higgs representations in an $SO(10)$ model, one may think that the amount of freedom one has for model building is enormous. However, in practice when one imposes the conditions that i) gauge coupling unification occurs, ii) that the intermediate scale is found to be below the GUT scale, and iii) that the GUT scale is high enough so that the proton lifetime exceeds current experimental bounds, only a handful of possible models survive [14, 15]. Furthermore, since any dark matter candidate must be part of a larger $SO(10)$ representation, that multiplet must be split, putting further constraints on the possible choice of field content.

In this letter, we choose one example of a scalar dark matter model with an intermediate scale gauge group given by $G_{\text{int}} = \text{SU}(3)_C \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L}$. We will examine the model labeled \mathbf{SA}_{3221} in [15] for which the dark matter is a scalar singlet originating in a **16** of $SO(10)$. In addition to SM fields, the model employs a **45** (or **210**) to break $SO(10)$ to G_{int} when the $(\mathbf{15}, \mathbf{1}, \mathbf{1})$ component (under $\text{SU}(4)_C \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R$) acquires a vacuum expectation value (vev). The intermediate scale gauge group is subsequently broken when the color singlet, right-handed triplet sitting in the **126** acquires a vev. All other components of the **126** are expected to have GUT scale masses. In addition to an explicit (GUT scale) mass term for the **16**, the scalar multiplet can have mass contributions from its couplings to the Higgs **45** and **126**. An explicit calculation of the fine-tuning needed to obtain a TeV scale mass for the singlet scalar dark matter candidate can be found in Appendix C of [15]. In the example given there, all members of the **16** are GUT scale except the scalar analog of e_R (\tilde{e}_R) which has an intermediate scale mass, and $\tilde{\nu}_R$ which has a weak scale mass.

Renormalization group evolution of the Higgs couplings and masses.—The renormalization group evolution between the weak scale and intermediate scale is almost identical to the SM. The only difference comes from the inclusion of the SM singlet dark matter candidate, $s \equiv \text{Re}[\tilde{\nu}_R]$. Below the intermediate scale, the scalar potential is relatively simple,

$$V_{\text{blw}} = \mu^2 |H|^2 + \frac{1}{2} \mu_s^2 s^2 + \frac{\lambda}{2} |H|^4 + \frac{\lambda_{sH}}{2} |H|^2 s^2 + \frac{\lambda_s}{4!} s^4. \quad (1)$$

In many ways, this resembles the minimal dark matter model often referred to as the Higgs portal [23, 24]. The mass of our dark matter candidate is given by $m_{\text{DM}}^2 = \lambda_{sH} v^2 / 2 + \mu_s^2$. Furthermore, fixing the dark matter mass will also fix λ_{sH} at the weak scale (taken here to be m_t) through the relic density (assuming standard thermal freeze-out): $m_{\text{DM}} \simeq 3.3 \lambda_{sH}$ TeV. In this paper, we compute the DM relic density using `micrOMEGAS` [25]. The evolution of the Higgs quartic coupling in the SM with and without the inclusion of the scalar s is shown in Fig. 1 by the green solid and dotted curves, respectively. The renormalization group equations (RGE) are run at the two-loop level¹ and one sees that the SM quartic coupling runs negative just above 10^{10} GeV [4] without the scalar contribution. With the scalar contribution, the running of λ would remain positive out to the GUT scale. Note that at the intermediate scale (determined by the conditions for gauge coupling unification; the running of the gauge couplings in \mathbf{SA}_{3221} is shown by thin black lines in Fig. 1), $M_{\text{int}} \simeq 10^9$ GeV, $\lambda > 0$. Gauge coupling unification also determines the GUT scale to be $M_{\text{GUT}} \simeq 1.5 \times 10^{16}$ GeV, which is high enough to evade the proton decay limit. Also shown is the running of λ_s (blue dash-dotted) and λ_{sH} (brown dashed).

Above the intermediate scale, it is necessary to include in addition to s , the right-handed doublet $\chi(\mathbf{1}, \mathbf{1}, \mathbf{2}, 1)$ which contains s , the Higgs triplet $\Delta(\mathbf{1}, \mathbf{1}, \mathbf{3}, 2)$ residing in the **126**, two heavy complex fields in addition to the SM Higgs doublet which all sit in a complex $\Phi(\mathbf{1}, \mathbf{2}, \mathbf{2}, 0)$, and finally the three right handed neutrinos sitting in the fermionic **16** matter representations. Above the intermediate scale, we write $\Phi = (\phi_1, \tilde{\phi}_2)$, $\tilde{\Phi} \equiv \sigma_2 \Phi^* \sigma_2$ (σ_a are the Pauli matrices), $\chi = (\chi^+, \chi^0)^T$, and

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}, \quad (2)$$

where $\phi_i = (\phi_i^0, \phi_i^-)^T$ is an $\text{SU}(2)_L$ doublet; $\tilde{\phi} \equiv i\sigma_2 \phi^*$.

¹ We use the three-loop RGEs for the top Yukawa and Higgs quartic couplings. We also include the two-loop electroweak threshold corrections according to Ref. [4]. We use the $\overline{\text{MS}}$ scheme up to the intermediate scale, and switch to the $\overline{\text{DR}}$ scheme at M_{int} .

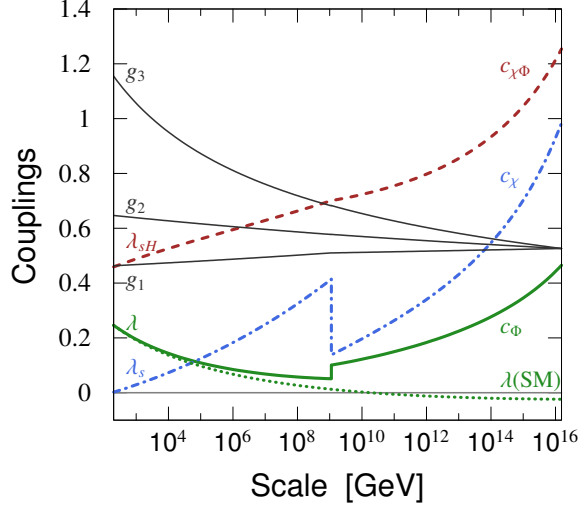


FIG. 1: Running of the quartic couplings of Higgs field, for selected inputs. The green solid, brown dashed, and blue dash-dotted lines show the running of λ , λ_{sH} , and λ_s , respectively, while the green dotted curve shows the running of λ in the SM. The gauge coupling curve running is also shown in thin black lines. Above the intermediate scale, the running of c_{Φ} , c_{χ} , and $c_{\chi\Phi}$ is shown using the matching conditions in (4). The free parameters are chosen as follows: At $Q = m_t$, $\lambda_s = 0$ and $\lambda_{sH} = 0.46$ (which corresponds to $m_{\text{DM}} \simeq 1.5$ TeV); At M_{int} , $\tilde{c}_{\Phi} = c'_{\Delta} = c_{\Phi\Delta} = c_{\chi\Delta} = c'_{\chi\Delta} = c'_{\chi\Phi} = 0$ and $c_{\Delta} = -c'_{\Phi\Delta} = -m_{\chi\Delta}/v_R = 0.05$. The non-zero couplings are taken so that the low-energy mass spectrum we consider here is realized.

Then a quartic potential can be written as

$$\begin{aligned}
 V_{\text{abv}}^{(4)} = & \frac{c_{\Delta}}{2} (\text{tr}(\Delta^{\dagger}\Delta))^2 + \frac{c'_{\Delta}}{4} \text{tr}(\Delta\Delta) \text{tr}(\Delta^{\dagger}\Delta^{\dagger}) \\
 & + \frac{c_{\Phi}}{2} (\text{tr}(\Phi^{\dagger}\Phi))^2 + \frac{\tilde{c}_{\Phi}}{4} \text{tr}(\tilde{\Phi}^{\dagger}\Phi) \text{tr}(\Phi^{\dagger}\tilde{\Phi}) \\
 & + c_{\Phi\Delta} \text{tr}(\Delta^{\dagger}\Delta) \text{tr}(\Phi^{\dagger}\Phi) + \frac{c_{\chi}}{2} |\chi|^4 + c_{\chi\Phi} |\chi|^2 \text{tr}(\Phi^{\dagger}\Phi) \\
 & + c_{\chi\Delta} |\chi|^2 \text{tr}(\Delta^{\dagger}\Delta) + c'_{\chi\Delta} \chi^{\dagger} [\Delta^{\dagger}, \Delta] \chi \\
 & + c'_{\Phi\Delta} \text{tr}(\Phi^{\dagger}\Phi [\Delta^{\dagger}, \Delta]) + c'_{\chi\Phi} \chi^{\dagger} \Phi^{\dagger} \Phi \chi + \dots \quad (3)
 \end{aligned}$$

Note that we have only included those quartic couplings which can be generated through RGE evolution, with the exception of the last two; $c'_{\Phi\Delta}$ is needed to split the masses of the two-Higgs doublet, Φ , while $c'_{\chi\Phi}$ is induced by the $c'_{\Phi\Delta}$ term via RGE effects.

The quartic terms that contain two powers of Δ , as well as the cubic coupling (see Eq. (5)) produce non-trivial tree-level threshold corrections at M_{int} , after Δ

acquires a vev and the heavy fields are integrated out:

$$\begin{aligned}
 \lambda &= c_{\Phi} - \frac{(c_{\Phi\Delta} + c'_{\Phi\Delta})^2}{c_{\Delta}}, \\
 \lambda_{sH} &= c_{\chi\Phi} - \frac{(c_{\Phi\Delta} + c'_{\Phi\Delta})[m_{\chi\Delta} + (c_{\chi\Delta} - c'_{\chi\Delta})v_R]}{c_{\Delta}v_R}, \\
 \lambda_s &= 3c_{\chi} - 3 \frac{[m_{\chi\Delta} + v_R(c_{\chi\Delta} - c'_{\chi\Delta})]^2}{c_{\Delta}v_R^2}, \quad (4)
 \end{aligned}$$

where $\langle\Delta\rangle = v_R T_-$ with $T_- \equiv (\sigma_1 - i\sigma_2)/2$. As is well known, these threshold effects always go in the direction of benefiting vacuum stability [7]. The evolution of the quartic couplings, c_{Φ} , c_{χ} , and $c_{\chi\Phi}$ above the intermediate scale are also shown in Fig. 1 using the matching conditions in (4). We use the one-loop RGEs for these quartic couplings. Although we do not explicitly display the running of all quartic terms above the intermediate scale, we have checked that although some run negative (notably c'_{Δ}), we have verified that the couplings satisfy sufficient conditions which guarantee stability of the vacuum up to the GUT scale.

The quadratic and cubic parts (which can lead to mass terms) of the potential can be written as

$$\begin{aligned}
 V_{\text{abv}}^{(2,3)} = & m_{\chi}^2 |\chi|^2 + m_{\Phi}^2 \text{tr}(\Phi^{\dagger}\Phi) + m_{\Delta}^2 \text{tr}(\Delta^{\dagger}\Delta) \\
 & + m_{\chi\Delta} (\tilde{\chi}^{\dagger} \Delta^{\dagger} \chi) + \text{h.c.}, \quad (5)
 \end{aligned}$$

where we take $m_{\chi\Delta}$ to be real for simplicity. The relevant matching conditions with the weak scale mass parameters are

$$\begin{aligned}
 \mu_s^2 &= m_{\chi}^2 + (c_{\chi\Delta} - c'_{\chi\Delta}) v_R^2 + 2m_{\chi\Delta} v_R, \\
 \mu^2 &= m_{\Phi}^2 + (c_{\Phi\Delta} + c'_{\Phi\Delta}) v_R^2, \quad (6)
 \end{aligned}$$

where the low energy fields are related to the high energy fields as $\phi_1 = H$ and $\chi^0 = (s + ia)/\sqrt{2}$.

The running of λ_s receives a large contribution from λ_{sH} , $d\lambda_s/d\ln Q = 12\lambda_{sH}^2/(4\pi)^2 + \dots$ and thus by demanding perturbativity of the couplings ($\lambda_i \lesssim 1/\beta_i$, where β_i is a relevant beta-function coefficient) up to the intermediate scale, we can set an upper bound on $\lambda_{sH} \lesssim 1.3$. However, requiring perturbativity of the c_i 's above the intermediate scale places a stronger bound on $\lambda_s(M_{\text{int}}) \lesssim 2.4$ which requires $\lambda_{sH}(m_t) \lesssim 0.9$. Non-zero values for other couplings further push the upper limit to $\lambda_{sH}(m_t) \lesssim 0.6$ in order to avoid singularities in the RGEs. Since λ_{sH} controls the annihilation cross section for s : $\sigma_{\text{ann}} v_{\text{rel}} \simeq \lambda_{sH}^2 / 16\pi m_{\text{DM}}^2$, and the relic density is proportional to $1/\langle\sigma_{\text{ann}} v_{\text{rel}}\rangle$, the upper limit on λ_{sH} corresponds to an upper limit to the DM mass $m_{\text{DM}} \lesssim 2$ TeV, similar to that in the minimal dark matter model [24] without an intermediate scale.

The Higgs mass parameter, μ^2 , must be negative in order to break the electroweak symmetry, and in the SM, μ^2 remains negative as it is run up to high energies. The

presence of the dark matter scalar however affects the running as $d\mu^2/d\ln Q = \lambda_{sH}\mu_s^2/(4\pi)^2 + \dots$ and causes μ^2 to run positive at higher renormalization scales [16]. In other words, the dark matter candidate can induce radiative electroweak symmetry similar to the mechanism in supersymmetric models [6]. As the running of μ depends on the combination $\lambda_{sH}\mu_s^2$ we can obtain a minimum value for μ_s (and hence m_{DM}) which is independent of the relic density constraint by maximizing λ_{sH} . We find that for $\lambda_{sH} = 0.6$, $\mu^2 > 0$ at the intermediate scale (at 1 TeV) when $\mu_s \gtrsim 360$ GeV (1150 GeV), corresponding to $m_{\text{DM}} \gtrsim 380$ GeV (1160 GeV). Here, we set $\lambda_s(m_t) = 0$. Taking the limits on λ_{sH} from the perturbativity of λ_s and the limit on μ_s from the requirement of radiative electroweak symmetry breaking, we find that the dark matter mass must lie in a restricted range (when demanding the more natural choice of symmetry breaking at 1 TeV) $m_{\text{DM}} = 1.2\text{--}2$ TeV.

When one imposes the constraint from the relic density, we obtain somewhat stronger bounds on λ_{sH} . In Fig. 2, we show the value of $\text{sgn}(\mu^2)|\mu|$ for $Q = M_{\text{int}}$ and 1 TeV as a function of $\lambda_{sH}(m_t)$. Here again, we set $\lambda_s(m_t) = 0$. As one can see that when $Q = M_{\text{int}}$, we have $\lambda_{sH}(m_t) > 0.2$ corresponding to $m_{\text{DM}} > 670$ TeV and when $Q = 1$ TeV, we have $\lambda_{sH}(m_t) > 0.41$ corresponding to $m_{\text{DM}} > 1.35$ TeV.

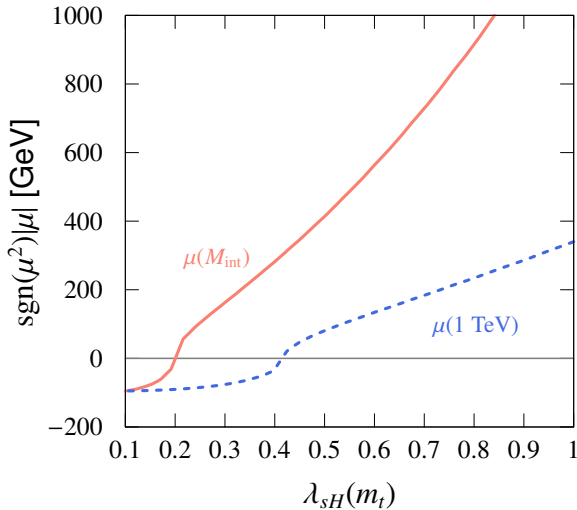


FIG. 2: The value of $\text{sgn}(\mu^2)|\mu|$ for $Q = M_{\text{int}}$ and 1 TeV as a function of $\lambda_{sH}(m_t)$. m_{DM} at the weak scale is determined from the requirement for the thermal relic abundance using $m_{\text{DM}} \approx 3.3\lambda_{sH}$ TeV.

The singlet DM candidate in our model can be probed in DM direct detection experiments. In Fig. 3, we show the spin-independent (SI) DM-nucleon scattering cross section σ_{SI} as a function of the DM mass. Here, we require the relic density condition to determine λ_{sH} . The lower solid (upper dashed) brown line shows the

result for which we use the nucleon matrix elements given in Ref. [26] (Ref. [27]). In either case, we obtain $\sigma_{\text{SI}} \simeq 10^{-45}$ cm². The gray shaded region is excluded by the current limit from the LUX experiment [28]. We also show the projected sensitivity of XENON1T [22] by the black dotted line. We find that all of the DM mass range can be probed at this experiment.

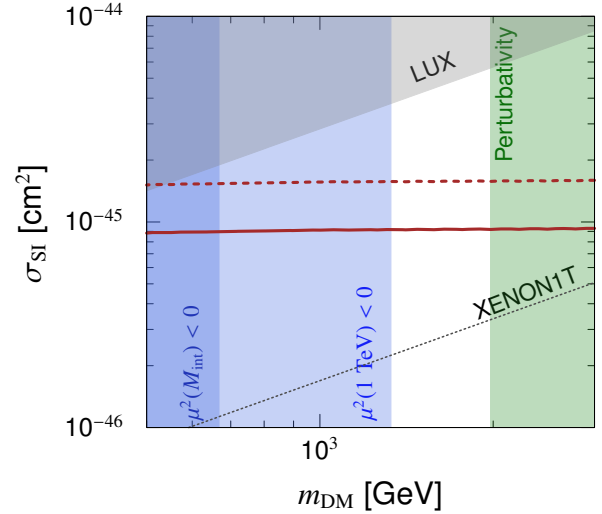


FIG. 3: The SI DM-nucleon scattering cross section as a function of m_{DM} . Here, λ_{sH} is determined from the relic density condition.

Summary.—We have presented an SO(10) model with gauge coupling unification made possible through an intermediate scale at $\simeq 10^9$ GeV. SO(10) is broken to $G_{\text{int}} = \text{SU}(3)_C \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L}$ when the right-handed triplet in the **126** obtains a vev. In this model, the lightest member of a complex scalar **16** is stable and plays the role of our dark matter candidate, s . The specific example discussed here can be viewed as a UV completion of the minimal (scalar) dark matter model. We have shown that in addition to gauge coupling unification, and a dark matter candidate, unlike the case in the SM, vacuum stability is achieved up to the GUT scale, and radiative electroweak symmetry breaking is triggered by the interactions of the dark matter and the SM Higgs. The latter result taken together with the requirement of perturbative couplings to the GUT scale limit the DM mass to lie between 1.35–2 TeV. This mass range should be probed in future direct detection experiments.

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