

A 125.5 GeV Higgs Boson in \mathcal{F} - $SU(5)$: Imminently Observable Proton Decay, A 130 GeV Gamma-ray Line, and SUSY Multijets & Light Stops at the LHC8

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We establish that the light Higgs boson mass in the context of the No-Scale Flipped $SU(5)$ GUT with TeV scale vector-like matter multiplets (flippons) is consistent with $m_h = 125.5 \pm 0.5$ GeV in the region of the best supersymmetry (SUSY) spectrum fit to low statistics data excesses observed by ATLAS in multijet and light stop 5 fb^{-1} SUSY searches at the LHC7. Simultaneous satisfaction of these disparate goals is achieved by employing a minor decrease in the $SU(5)$ partial unification scale M_{32} to lower the flippon mass, inducing a larger Higgs boson mass shift from the flippon loops. The reduction in M_{32} , which is facilitated by a phenomenologically favorable reduction of the low-energy strong coupling constant, moreover suggests an imminently observable $(e|\mu)^+\pi^0$ proton decay with a central value time scale of 1.7×10^{34} years. At the same point in the model space, we find a lightest neutralino mass of $m_\chi = 145$ GeV, which is suitable for the production of 130 GeV monochromatic gamma-rays through annihilations yielding associated Z -bosons; a signal with this energy signature has been identified within observations of the galactic center by the FERMI-LAT Space Telescope. In conjunction with direct correlations to the fate of the ATLAS multijet and light stop production channels presently being tested at the LHC8, we suggest that the reality of a 125.5 GeV Higgs boson affords a particularly rich company of specific and imminently testable associated observables.

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I. INTRODUCTION

The recent combined discovery of the CP-even light Higgs boson around $m_h = 125\text{--}126$ GeV by the ATLAS [1], CMS [2], and CDF/D0 [3] Collaborations has sparked detailed examinations into what regions of the SUSY parameter space remain viable, and which are very constrained, if not outright excluded. Motivated by these new precision experimental measurements, we endeavor here to commensurately enhance the precision of prior theoretical estimates [4, 5] for the range of Higgs boson masses consistent with the No-Scale Flipped $SU(5)$ model with vector-like matter (flippons) [4–23], dubbed \mathcal{F} - $SU(5)$ for short. The mutual consistency of a 125–126 GeV Higgs boson mass with an explanation for the tantalizing positive excesses [20, 21] observed at low statistics by ATLAS [24, 25] and CMS [26] in the $\sqrt{s} = 7$ TeV SUSY search dramatically narrows the \mathcal{F} - $SU(5)$ model space, with interesting and specific implications for ongoing collider, indirect dark matter detection, and proton decay experiments.

The central mass peak of the new particle observed at the LHC exists at 125.3 GeV and 126.0 GeV, according respectively to the CMS [2] and ATLAS [1] collaborations. In each case, statistical and systematic errors of about half a GeV are suggested. Lacking an official combination of the two experiments' statistics, a more broad treatment might follow something akin to the proposal

$m_h = 125.0 \pm 1.0$ (exp) ± 1.5 (theory) GeV of Ref. [27], noting in particular that our capacity to computationally model the predicted Higgs value is likewise imprecise. However, it is our interest in the present work to ascertain the restrictions that a Higgs mass very close to the average of centrally reported values would place on an otherwise successful model. Inasmuch as this mass is somewhat heavier than what may be comfortably achieved in typical GUT models without resorting to very heavy sparticles, we are interested in demonstrating a counterexample satisfying that vital emerging constraint which does not resort to taking extremities in the admissible lower bound. In particular, we presently investigate the possibility that a large portion of the burden could be absorbed (in conjunction with the described flippon loops) by a lowering of the strong coupling α_s at M_Z , an accommodation to which the flipped $SU(5)$ GUT is particularly well historically adapted [28]. The Higgs mass constraint $m_h = 125.5 \pm 0.5$ adopted for this work is thus purposefully rather strict.

II. \mathcal{F} - $SU(5)$: THE MODEL

No-Scale \mathcal{F} - $SU(5)$ is built upon the foundation of the Flipped $SU(5)$ [29–31] Grand Unified Theory (GUT), two pairs of hypothetical TeV-scale flippon multiplets of mass M_V derived from local F-Theory model build-

ing [32–36], and the dynamically established boundary conditions of No-Scale supergravity [37–41]. In flipped $SU(5) \times U(1)_X$ models, there are two unification scales: the $SU(3)_C \times SU(2)_L$ unification scale M_{32} and the $SU(5) \times U(1)_X$ unification scale $M_{\mathcal{F}}$. To separate the M_{32} and $M_{\mathcal{F}}$ scales and obtain true string-scale gauge coupling unification in free fermionic string models [32, 42] or the decoupling scenario in F-theory models [33, 34], we introduce vector-like particles, called flippons, which form complete flipped $SU(5) \times U(1)_X$ multiplets. In the most elementary No-Scale scenario, $M_0 = A = B_\mu = 0$ at the ultimate unification scale $M_{\mathcal{F}}$, while the entire group of low energy SUSY breaking soft-terms evolve down from just one non-zero parameter $M_{1/2}$. As a result, the particle spectrum is proportional to $M_{1/2}$ at leading order, rendering the bulk “internal” physical properties invariant under an overall rescaling. The matching condition between the low-energy value of B_μ that is demanded by electroweak symmetry breaking (EWSB) and the high-energy $B_\mu = 0$ boundary condition is quite difficult to reconcile under the renormalization group equation (RGE) running. The solution at hand applies modifications to the β -function coefficients that are produced by the flippon loops. Naturalness in view of the gauge hierarchy and μ problems suggest that the flippon mass M_V should be of the TeV order. Avoiding a Landau pole for the strong coupling constant restricts the set of vector-like flippon multiplets which may be given a mass in this range to only two constructions with flipped charge assignments, which have been explicitly realized in the F-theory model building context [32–34]. In either case, the (formerly negative) one-loop β -function coefficient of the strong coupling α_3 becomes precisely zero from M_{32} to M_V , flattening the RGE running of the gaugino mass M_3 and the strong coupling α_3 between these two scales. The wide residual gap which this leaves between the α_5 and α_X couplings at M_{32} facilitates a quite important secondary running phase up to the final $SU(5) \times U(1)_X$ unification scale, which may be elevated by 2-3 orders of magnitude into adjacency with the Planck mass, where the $B_\mu = 0$ boundary condition fits like hand to glove [6, 43, 44].

III. HIGGS BOSON MASS SHIFT

The correlation of the predicted \mathcal{F} - $SU(5)$ Higgs boson mass in Ref. [5] with the collider measured value was achieved via contributions to the lightest Higgs boson from the vector-like flippons, computed from the RGE improved one-loop effective Higgs potential approach [45, 46], resulting in a 3-4 GeV upward shift in the Higgs boson mass to the experimentally measured range. The relevant shift possesses a leading dependence on the flippon mass M_V , with larger shifts corresponding to lighter vector-like flippons. The recently improved measurements that further constrain the Higgs boson mass to within a more narrow corridor of $m_h = 125.5 \pm 0.5$

GeV necessitate a refinement in our calculations from our previous \mathcal{F} - $SU(5)$ estimate of $m_h \simeq 124$ GeV [20, 23] at that location in the model space capable of explaining the multijet and light stop excesses. In order to accomplish such, the flippon mass M_V must be lowered further to accommodate a larger flippon contribution to the Higgs boson mass.

A multi-axis χ^2 test was applied to those ATLAS and CMS searches exhibiting the largest production of events beyond the data-driven Standard Model predictions within the high-energy framework of \mathcal{F} - $SU(5)$ [20, 21]. It was shown that the best fit SUSY mass to these searches occurred at $M_{1/2} \simeq 700$ GeV. The No-Scale \mathcal{F} - $SU(5)$ SUSY spectrum possesses the rather distinct characteristic of leading order *en masse* proportionality to only the single dimensionful parameter $M_{1/2}$. In essence, the internal physics of \mathcal{F} - $SU(5)$ are largely invariant under a numerical rescaling of only $M_{1/2}$. For this reason, we must leave $M_{1/2}$, which is defined at the ultimate $SU(5) \times U(1)_X$ unification scale $M_{\mathcal{F}}$, relatively unchanged, while simultaneously shifting the flippon mass M_V , since the event landscape at the LHC is chiefly correlated to $M_{1/2}$. This rescaling of the mass scale $M_{1/2}$ shares an analogous historical bond with the fixing of the Bohr atomic radius $a_0 = 1/(m_e\alpha)$ in terms of the physical electron mass and charge, through minimization of the electron energy [47]. In each of these two instances, the spectrum scales according to variation in the selected constants, while leaving the relative internal structure of the model intact.

The key to realizing a modest shift in M_V while maintaining stability in $M_{1/2}$ at the ultimate unification scale $M_{\mathcal{F}}$ is the recognition that some flexibility remains with regards to the positioning of the penultimate $SU(3)_C \times SU(2)_L$ unification scale M_{32} . By subjecting M_{32} to minor variations, we can deviate the flippon mass M_V through the RGE running, while preserving reasonable stability of $M_{1/2}$ at $M_{\mathcal{F}}$. In practice, small transitions in the M_{32} scale are executed in reverse by programmatic variation of the low-energy strong coupling constant α_s . The strong coupling constant fulfills a low-energy boundary condition for the RGE running, thus small movements at low-energy, while keeping $M_{1/2}$ at $M_{\mathcal{F}}$ reasonably fixed, translate into minor shifts in M_{32} . These small variations in the low and high-energy boundary conditions thus map onto the desired motion of the flippon mass M_V , while maintaining the favorable phenomenology dictated by a particular $M_{1/2}$.

To quantify the allowable latitude of variation for the low-energy strong coupling constant α_s , we turn to the most recent report of precision electroweak scale measurements of this parameter [48], which announces the refined result $\alpha_s = 0.1161_{-0.0048}^{+0.0041}$. Considering that our previous work in its entirety [4–23] employed a value of $\alpha_s = 0.1172$, quite soundly within the measured QCD uncertainties, though a bit above the centrally measured value, we may presently take advantage of an opportunity to displace this coupling downward, toward and

through the central value. In fact, shifts in α_s and M_{32} operate with a parallel directionality, inducing a flow of the flippon mass M_V that is likewise of the same sign. Hence, the small shortfall between our previous Higgs boson mass calculation of $m_h \simeq 124$ GeV and the now more highly constrained Higgs boson mass measurements centered at 125.5 GeV may be attributed to the application of an overly large strong coupling constant α_s at the low-energy boundary condition.

As we shall explicitly show, lowering α_s to and slightly below the recently reported central value does indeed permit a sufficient decrease in the flippon mass M_V to allow an $m_h \simeq 125.5$ GeV Higgs boson mass at the well of our cumulative χ^2 best fit to the multijet and light stop excesses at the LHC. Moreover, oscillating the strong coupling constant α_s around its central empirically derived value in turn oscillates the \mathcal{F} - $SU(5)$ computed Higgs boson mass around $m_h \sim 125$ GeV, prominently emphasizing the natural and non-trivial \mathcal{F} - $SU(5)$ SUSY GUT correlation between two independent high-precision low-energy experimental measurements, specifically the strong coupling constant α_s and light Higgs boson mass m_h . In Section (VI), we will supplement numerically based observations on the co-dependency of M_{32} and α_s with an analytical treatment emphasizing the vital role that lowering the GUT scale M_{32} plays in hastening another key indicator of post Standard Model physics: the decay of the proton.

IV. HIGGS BOSON MASS CALCULATIONS

We vary the strong coupling constant α_s to ascertain those regions of the model space capable of handily generating an $m_h = 125.5 \pm 0.5$ GeV Higgs boson mass. To first constrain the model space to only those regions satisfying current experiment, we apply the set of stable “bare-minimal” experimental constraints of Ref. [12], consisting of the top quark mass $172.2 \text{ GeV} \leq m_t \leq 174.4 \text{ GeV}$, 7-year WMAP cold dark matter relic density $0.1088 \leq \Omega_{\text{CDM}} h^2 \leq 0.1158$ [49], and precision LEP constraints on the SUSY mass content. In the bulk of the surviving model space the lightness of the stau, which is itself a potential future target for direct collider probes by the forthcoming $\sqrt{s} = 14$ TeV LHC, is leveraged to facilitate an appropriate dark matter relic density via stau-neutralino coannihilation. The SUSY particle masses and relic densities are calculated with `MicrOMEGAs 2.4` [50], via application of a proprietary modification of the `SuSpect 2.34` [51] codebase to evolve the flippon-enhanced RGEs. The resultant model space is then used as a basis to compute the flippon contribution to the Higgs boson mass per Ref. [46], as a first approximation in deriving the total Higgs boson mass m_h . This carves out a very narrow strip of model space as a function of $(M_{1/2}, M_V)$ that will shift as a unit along the M_V axis in response to small variations in α_s . These narrow strips of model space are subsequently utilized to

complete more detailed calculations of the Higgs boson mass to be described shortly. We find an $M_{32} \simeq 7 \times 10^{15}$ GeV and $\alpha_s = 0.1145$ to be ideal for achieving a Higgs boson mass of $m_h \sim 125.5$ GeV.

A brief comparison is in order between the structure of the vector-like matter content appearing in the models of Ref. [46] and that employed in \mathcal{F} - $SU(5)$. The \mathcal{F} - $SU(5)$ construction adopts the multiplets

$$\begin{aligned} XF &= (\mathbf{10}, \mathbf{1}) \equiv (XQ, XD^c, XN^c) ; \overline{XF} = (\overline{\mathbf{10}}, -\mathbf{1}) \\ Xl &= (\mathbf{1}, -\mathbf{5}) ; \overline{Xl} = (\mathbf{1}, \mathbf{5}) \equiv XE^c, \end{aligned} \quad (1)$$

where XQ , XD^c , XE^c and XN^c carry the same quantum numbers as the quark doublet, right-handed down-type quark, charged lepton and neutrino, respectively. This is one of two possibilities realizing complete $SU(5) \times U(1)_X$ multiplets at the TeV scale that avoids a Landau pole for the strong coupling constant; alternatively, the pair of $SU(5)$ singlets may be discarded, but phenomenological consistency then requires the substantial application of unspecified GUT thresholds. Both scenarios have been explicitly realized in the F -theory and free-fermionic model building contexts [32–34]. The $SU(3)_C \times SU(2)_L$ gauge symmetry unification is not consequently shifted far from the traditional GUT scale around 10^{16} GeV, although the vector-like multiplets do elevate the $SU(5) \times U(1)_X$ gauge symmetry unification into close proximity with the string scale.

In Ref. [46], the author has considered vector-like matter corresponding to the charge assignments

$$(XQ, XQ^c) + (XD, XD^c) + 2(XE, XE^c), \quad (2)$$

which do not form complete $SU(5)$ multiplets, but which nevertheless do preserve a traditional GUT coupling unification scale. The detailed string-theoretic construction of GUTs with such vector-like particles remains an open question, so this model is classified as phenomenological. Nevertheless, the Yukawa superpotential terms

$$XQ^c XDH_u \quad \& \quad XQXD^c H_d \quad (3)$$

employed in Ref. [46] to increase the lightest CP-even Higgs boson mass are directly applicable to the \mathcal{F} - $SU(5)$ field content, such that the discussions and formulae established in that work are likewise directly relevant to our own analysis.

The mechanism for generation of the flippon contributions to the Higgs boson mass is the following pair of Yukawa interaction terms between the MSSM Higgs boson and the vector-like flippons in the superpotential.

$$W = \frac{1}{2} Y_{xd} X F X F h + \frac{1}{2} Y_{xu} \overline{XF} \overline{XF} \overline{h} \quad (4)$$

Being vector-like rather than chiral, the flippons are also afforded a proper “diagonal” Dirac mass. After the $SU(5) \times U(1)_X$ gauge symmetry is broken down to the Standard Model, the relevant Yukawa couplings are

$$W = Y_{xd} XQXD^c H_d + Y_{xu} XQ^c XDH_u . \quad (5)$$

We employ the RGE improved one-loop effective Higgs potential approach to calculate the contributions to the lightest CP-even Higgs boson mass from the vector-like particles [45, 46]. The relevant shift in the Higgs boson mass-square is approximated as [52]

$$\begin{aligned} \Delta m_h^2 &= -\frac{N_c M_Z^2}{8\pi^2} \times \cos^2 2\beta (\hat{Y}_{xu}^2 + \hat{Y}_{xd}^2) t_V \\ &+ \frac{N_c v^2}{4\pi^2} \times \hat{Y}_{xu}^4 (t_V + \frac{1}{2} X_{xu}), \end{aligned} \quad (6)$$

with

$$\begin{aligned} \hat{Y}_{xu} &= Y_{xu} \sin \beta \quad ; \quad \hat{Y}_{xd} = Y_{xd} \cos \beta \\ \tilde{A}_{xu} &= A_{xu} - \mu \cot \beta \quad ; \quad t_V = \ln \frac{M_S^2 + M_V^2}{M_V^2} \\ X_{xu} &= \frac{-2M_S^2(5M_S^2 + 4M_V^2) + 4(3M_S^2 - 2M_V^2)\tilde{A}_{xu}^2 + \tilde{A}_{xu}^4}{6(M_V^2 + M_S^2)^2}, \end{aligned} \quad (7)$$

where N_c is the number of colors carried by the vector-like fields, M_S is the geometric mean of the left-handed and right-handed stop masses at low energy, and A_{xu} is the soft SUSY breaking trilinear term for the Yukawa superpotential element $Y_{xu} X Q^c X D H_u$.

The Yukawa couplings Y_{xu} and Y_{xd} are required to be smaller than about 3 throughout the full running up to the unification scale due to the perturbative bound. Thus, they are smaller than about 1.02 at low energies, as detailed in Ref. [52]. Specifically, from the two-loop RGE analyses, we determined that the maximal Yukawa couplings Y_{xu} are about 0.96, 1.03, and 1.0 for $\tan \beta = 2$, $\tan \beta \sim 23$, and $\tan \beta = 50$, respectively [52], and thus choose a working value of $Y_{xu} = 1.0$. Note, in particular, that the Yukawa couplings Y_{xu} and Y_{xd} remain fully distinct, even in the flipped $SU(5) \times U(1)_X$ models. Thus, for simplicity, we here assume that $Y_{xd} = 0$. With both Y_{xu} and Y_{xd} non-zero, one may study the RGE running and Higgs boson mass similarly, and the established results for the Higgs boson mass are found to broadly hold [52]. The corrected Higgs boson mass will be

$$m_h = \sqrt{(\tilde{m}_h)^2 + \Delta m_h^2}, \quad (8)$$

where \tilde{m}_h is the Higgs boson mass in the Minimal Supersymmetric Standard Model (MSSM). We implement this detailed Higgs boson mass computational methodology on points sampled from the residual strip of model space surviving subsequent to application of the ‘‘bare-minimal’’ constraints [12] and the simplified first-level Higgs boson mass approximation detailed in Ref. [5]. In that preliminary analysis, a formula for the Higgs mass-square shift Δm_h^2 attributable to the vector-like flippon multiplets (and also more weakly to the soft term mass approximated by $M_S \simeq 2M_{1/2}$) is adopted that implements the leading mass scale dependencies identified in Section III of Ref. [46]. This allows us to reasonably extrapolate a detailed computation of the corrected Higgs mass across the \mathcal{F} - $SU(5)$ model space, isolating a narrow region of the model space supporting $123 \lesssim m_h \lesssim 127$

GeV, as illustrated in Figure (1).

$$\begin{aligned} m_h &\simeq \sqrt{\tilde{m}_h^2 + (87.81 \text{ GeV})^2 \times \left(\ln x - \frac{5}{6} + \frac{1}{x} - \frac{1}{6x^2} \right)} \\ x &\equiv 1 + \left(\frac{2M_{1/2}}{M_V} \right)^2 \end{aligned} \quad (9)$$

Larger shifts are seen to correspond to lighter vector-like multiplets. This effect works in tandem with the top quark mass, which varies inversely with M_V in the viable \mathcal{F} - $SU(5)$ model space [5], and whose elevation likewise lifts the bare Higgs mass prediction \tilde{m}_h . For both reasons, larger net values of the Higgs mass m_h occur toward the lower \mathcal{F} - $SU(5)$ model boundary for M_V , just prior to the extremity of a single standard deviation from the top quark world average.

Note that no 5 fb^{-1} 7 TeV LHC constraints are imposed upon the strip in Figure (1). The enduring region is comprised of a narrow strip of space confined to $400 \lesssim M_{1/2} \lesssim 900 \text{ GeV}$, $19.4 \lesssim \tan \beta \lesssim 23$, and $1000 \lesssim M_V \lesssim 6500 \text{ GeV}$. The border at the minimum $M_{1/2} \simeq 400 \text{ GeV}$ is required by the LEP constraints, while the maximum boundary at $M_{1/2} \simeq 900 \text{ GeV}$ prevents a charged stau LSP. We superimpose upon the plot space in Figure (1) the cumulative χ^2 curve for a set of ATLAS SUSY searches exhibiting appreciable signal significance over the expected background in accordance with the procedures detailed in Refs. [20, 21]. The reduction in the strong coupling constant to $\alpha_s = 0.1145$ is carried over into this analysis. Figure (1) depicts the successful preservation of a best SUSY mass fit at $M_{1/2} \simeq 700 \text{ GeV}$, in tandem with a smaller $M_V \simeq 2600 \text{ GeV}$, which provides the desired upward shift in the Higgs boson mass to $m_h \simeq 125.5 \text{ GeV}$. Accordingly, the SUSY mass spectrum of the benchmark model of Table I at the χ^2 well remains rather similar to that of Ref. [20], with a bino-dominated LSP $m_{\tilde{\chi}_1^0} = 145 \text{ GeV}$, light stop $m_{\tilde{t}_1} = 777 \text{ GeV}$, gluino $m_{\tilde{g}} = 945 \text{ GeV}$, and heavy squark $m_{\tilde{u}_L} = 1489 \text{ GeV}$. We also provide a somewhat heavier spectrum in Table II, as graphically annotated in Figure (1), that is within one standard deviation of the χ^2 minimum, with a bino-dominated LSP $m_{\tilde{\chi}_1^0} = 181 \text{ GeV}$, light stop $m_{\tilde{t}_1} = 937 \text{ GeV}$, gluino $m_{\tilde{g}} = 1133 \text{ GeV}$, and heavy squark $m_{\tilde{u}_L} = 1739 \text{ GeV}$.

TABLE I: Spectrum (in GeV) for $M_{1/2} = 708 \text{ GeV}$, $M_V = 2612 \text{ GeV}$, $m_t = 174.4 \text{ GeV}$, $\tan \beta = 21.83$. Here, $\Omega_{\text{CDM}} h^2 = 0.1110$ and the lightest neutralino is greater than 99% Bino.

$\tilde{\chi}_1^0$	145	$\tilde{\chi}_1^\pm$	311	\tilde{e}_R	264	\tilde{t}_1	777	\tilde{u}_R	1367	m_h	125.3
$\tilde{\chi}_2^0$	311	$\tilde{\chi}_2^\pm$	1134	\tilde{e}_L	750	\tilde{t}_2	1259	\tilde{u}_L	1489	$m_{A,H}$	1230
$\tilde{\chi}_3^0$	1131	$\tilde{\nu}_{e/\mu}$	745	$\tilde{\tau}_1$	153	\tilde{b}_1	1224	\tilde{d}_R	1420	m_{H^\pm}	1233
$\tilde{\chi}_4^0$	1133	$\tilde{\nu}_\tau$	726	$\tilde{\tau}_2$	733	\tilde{b}_2	1348	\tilde{d}_L	1491	\tilde{g}	945

The cumulative distribution function (CDF) percentage labeled on the right-hand axis of the χ^2 plot in Figure (1) is a statistical tool that establishes the fraction

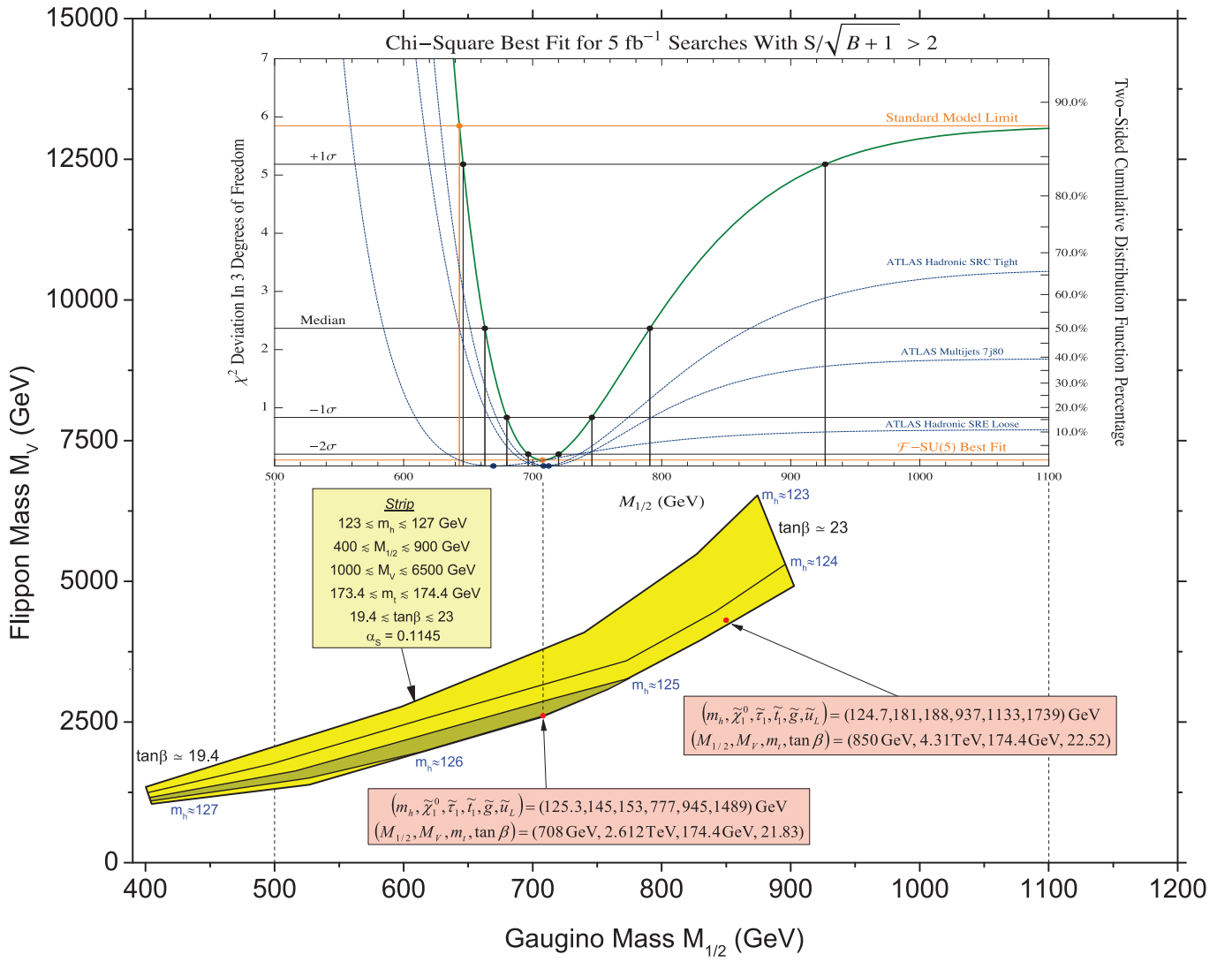


FIG. 1: We depict the experimentally viable parameter space of No-Scale \mathcal{F} -SU(5) as a function of the gaugino mass $M_{1/2}$ and flippon mass M_V . The surviving model space after application of the bare-minimal constraints of Ref. [12] and Higgs boson mass calculations is illustrated by the yellow strip, representing $123 \lesssim m_h \lesssim 127 \text{ GeV}$. A visual indication of the freedom that may be realized by relaxing the strong bounds on the Higgs mass is given by inclusion of outer contours at 123 and 124 GeV, corresponding to a weaker interpretation of the Higgs mass central value and computational uncertainties, and an outer contour at 127 GeV representing a stronger interpretation of the central value. The inset diagram (with linked horizontal axis) is the multi-axis cumulative χ^2 fitting of Ref. [20], though here using a slightly altered strong coupling constant of $\alpha_s = 0.1145$. Additionally, updates to certain of the relevant published [24] background estimates have actually enhanced the agreement in the mass scale favored across the displayed search channels. Highlighted are benchmarks at $M_{1/2} = 708 \text{ GeV}$ and $M_{1/2} = 850 \text{ GeV}$; the complete SUSY mass spectra are given in Tables I and II.

of trials (for a fixed number of statistically independent variables) where Gaussian fluctuation of each of variable will yield a net deviation from the null hypothesis that is not larger than the corresponding χ^2 value referenced on the left-hand axis. The two-sided limit as adopted for the χ^2 best-fit minimization against $M_{1/2}$ using searches with visible excesses is appropriate when meaningful deviations may be anticipated in either direction away from the median. Particularly small χ^2 values imply a better fit to the data by the assumed signal (at a certain confidence level) than what could be attributed to random

fluctuations around the SM, while excessively large χ^2 values disfavor a given $M_{1/2}$ relative to the SM-only null hypothesis. In this case, the usual 1 and 2- σ 68% and 95% consistency integrations enclose areas symmetrically distributed about the mean, such that centrally inclusive boundary lines are drawn at the CDF percentages 2.2%, 15.9%, 84.1%, and 97.7%. It should be noted that this method implicitly assumes the displayed channels to be statistically uncorrelated, which is not perfectly applicable in the current case. A compensating reduction in the effective degrees of freedom from the nominal value

TABLE II: Spectrum (in GeV) for $M_{1/2} = 850$ GeV, $M_V = 4310$ GeV, $m_t = 174.4$ GeV, $\tan\beta = 22.52$. Here, $\Omega_{\text{CDM}}h^2 = 0.1129$ and the lightest neutralino is greater than 99% Bino.

$\tilde{\chi}_1^0$	181	$\tilde{\chi}_1^\pm$	382	\tilde{e}_R	315	\tilde{t}_1	937	\tilde{u}_R	1596	m_h	124.7
$\tilde{\chi}_2^0$	382	$\tilde{\chi}_2^\pm$	1301	\tilde{e}_L	875	\tilde{t}_2	1464	\tilde{u}_L	1739	$m_{A,H}$	1413
$\tilde{\chi}_3^0$	1298	$\tilde{\nu}_{e/\mu}$	872	$\tilde{\tau}_1$	188	\tilde{b}_1	1436	\tilde{d}_R	1655	m_{H^\pm}	1416
$\tilde{\chi}_4^0$	1300	$\tilde{\nu}_\tau$	848	$\tilde{\tau}_2$	854	\tilde{b}_2	1568	\tilde{d}_L	1740	\tilde{g}	1133

of three would have the effect of marginally lowering the quoted CDF scale values relative to the left-hand χ^2 axis, slightly compressing the displayed error margins.

V. MULTIJETS AND LIGHT STOPS

It is not, in our opinion, coincidental that the only search strategies to exhibit appreciable signal strength thus far, amongst the still relatively small 5 fb^{-1} of accumulated statistics, consist of multijets. In fact, this feature is handily in accordance with our expectations for No-Scale \mathcal{F} - $SU(5)$, as outlined in Refs. [10, 11]. The model preference for multijet events results from leveraging the flippon multiplets to facilitate the flat RGE running from M_{32} to M_V , engendering the distinctive mass texture $M(\tilde{t}_1) < M(\tilde{g}) < M(\tilde{q})$, featuring a light stop and gluino that are both substantially lighter than all other squarks. The spectrum produces a characteristic event topology starting with the pair production of heavy first or second generation squarks \tilde{q} and/or gluinos \tilde{g} in the initial hard scattering process, with each heavy squark likely to yield a quark-gluino pair $\tilde{q} \rightarrow q\tilde{g}$ in the cascade decay. Each gluino will most likely decay to light stops \tilde{t}_1 via $\tilde{g} \rightarrow \tilde{t}_1\tilde{t} \rightarrow t\tilde{t}\tilde{\chi}_1^0$, or absent light stops, $\tilde{g} \rightarrow q\tilde{q}\tilde{\chi}_1^0$, with each gluino producing 2–6 jets, where the gluino-mediated stop channel produces the maximum of six jets.

In Refs. [10, 11], we advocated for a rather soft cut on hadronic jets of $p_T > 20$ GeV, in order to capture as many large multijet events with at least nine jets as possible. While we are certainly pleased that ATLAS has focused attention on multijet events, as evidenced by the powerful search strategies highlighted in Figure (1), the p_T cut on jets in these strategies is too hard to preserve events with at least nine jets in only 5 fb^{-1} of 7 TeV statistics. Consequently, we find the very recent 5 fb^{-1} light stop search results in ATLAS Ref. [53] quite exciting, where a softer jet cut of $p_T > 25$ GeV was employed, producing a notable excess of pair-produced light stop event candidates in the E_T^{Miss} bin correlated to about twice the $m_\chi = 145$ GeV lightest neutralino mass [14] of the best fit benchmark from Table I. A close scrutiny of the ATLAS light stop search reveals Figure (5) of Ref. [53], clearly illustrating that all of the 2.5σ excess events emanated from events with 9–10 jets in both the electron and muon channels, entirely consistent with our SUSY discovery prescription of Refs. [10, 11].

We emphasize that the magnitude of the observed excess [19, 20] and the value of the missing energy at which any observed excess is maximized [14] are independent discriminants in the \mathcal{F} - $SU(5)$ context, which must each ultimately link fundamental SUSY processes back to a single consistent gaugino mass $M_{1/2}$, the wellspring from which all correlations within this model flow.

A full-scale analysis of the ATLAS light stop search of Ref. [53] in the context of an \mathcal{F} - $SU(5)$ framework is in progress [54], and indeed we have already established that the best SUSY mass fit of the light stop excesses are well correlated with the three multijet ATLAS searches depicted in Figure (1). Moreover, since the search strategy of Ref. [53] targets pair-produced light stops, while the other three searches of Figure (1) target pair-produced gluinos and/or squarks, this light stop search should be statistically independent of these other high-productivity searches. The SRC Tight (4 jets and $p_T > 60$ GeV for j2–j4), SRE Loose (6 jets and $p_T > 60$ GeV for j2–j4, $p_T > 40$ GeV for j5–j6), and 7j80 (≥ 7 jets and $p_T > 80$ GeV) searches are also statistically independent as a result of the larger $p_T > 80$ GeV cut for 7j80, such that to capture the same event between more than one of these three searches would require lowering the number of jets in 7j80 events to below seven, which would in turn necessitate raising the p_T cut to greater than 80 GeV, though the jet cut is lower than 80 GeV for SRC Tight and SRE Loose.

VI. FAST PROTON DECAY

Instability of the proton is an essential signature of GUTs, the merger of the Standard Model forces necessarily linking quarks to leptons, and providing a narrow channel $p \rightarrow (e|\mu)^+\pi^0$ of dimension six decay via heavy gauge boson exchange. The 50-kiloton (kt) water Čerenkov detector of the Super-Kamiokande facility has now set a lower bound of 1.4×10^{34} years for the partial lifetime in the $p \rightarrow e^+\pi^0$ mode, and somewhat less for $\mu^+\pi^0$ [55]. The predicted lifetime for each of these two channels may be comparable in flipped $SU(5)$. However, this construction evades the dangerously rapid $p \rightarrow K^+\bar{\nu}$ dimension five triplet Higgsino mediated decay that plagues other SUSY GUTs by way of the missing-partner mechanism [31], which naturally facilitates Higgs doublet-triplet splitting without the side effect of strong triplet mixing.

We have previously [35] provided a comprehensive dictionary for the translation in closed form of low energy experimental observables into their flipped $SU(5)$ counterparts expressed at the high energy scale, accurate up to leading effects in the second order. In particular, the two couplings α_5 and α_X of the flipped $SU(5) \times U(1)_X$ gauge group and the GUT scale mass M_{32} — at which the partially unified $SU(5)$ sector decomposes into $SU(3) \times SU(2)_L$ and the Abelian phase associated with this breaking of rank from 4 to 3 likewise

remixes with the quarantined $U(1)_X$ subgroup to form the weak scale hypercharge $U(1)_Y$ — may be written in terms of the Z-boson mass M_Z , the values of the strong and electromagnetic couplings (α_s, α_{em}) and Weinberg angle $\sin^2 \theta_W$ at M_Z , the top quark mass m_t and the low-energy SUSY particle spectrum (for computation of threshold corrections), and the β -function renormalization group coefficients at the first and second loops. Our current interest is a direct isolation of dependencies in M_{32} and $1/\alpha_5$, to which the proton lifetime τ_p in the $p \rightarrow (e|\mu)^+ \pi^0$ dimension six modes is proportional in the fourth and second powers respectively, upon variation of the strong coupling $\alpha_s(M_Z)$.

To facilitate this effort, we borrow the following from Eq. Set (69) of Ref. [35].

$$\begin{aligned} M_{32} &= M_Z \times \exp \left\{ \frac{2\pi(\alpha_s \Theta_W - \alpha_{em} \Xi_3)}{\alpha_{em} \alpha_s (b_2 - b_3)} \right\} \\ \frac{1}{\alpha_5} &= \frac{\alpha_{em} \Xi_3 b_2 - \alpha_s \Theta_W b_3}{\alpha_{em} \alpha_s (b_2 - b_3)} \end{aligned} \quad (10)$$

In the prior, $\Theta_W \equiv \sin^2 \theta_W + \alpha_{em} (\xi_2 - \zeta_2)/2\pi$ and $\Xi_3 \equiv 1 + \alpha_s (\xi_3 - \zeta_3)/2\pi$, where ξ_i and ζ_i respectively incorporate effects associated to the i^{th} gauge coupling of mass threshold crossings and the second quantum loop. The b_i parameters are the single loop β -function coefficients for the running of the i^{th} coupling. The α_i to which we refer are the usual GUT-normalized $SU(3) \times SU(2)_L \times U(1)_Y$ couplings, with boundaries expressed as follows.

$$\begin{aligned} \alpha_1(M_Z) &= \frac{5 \alpha_{em}}{3 (1 - \sin^2 \theta_W)} \\ \alpha_2(M_Z) &= \frac{\alpha_{em}}{\sin^2 \theta_W} \\ \alpha_3(M_Z) &= \alpha_s \end{aligned} \quad (11)$$

The crucial observation is that Θ_W has no leading dependence on α_s , while Ξ_3 represents a deviation from unity that is linearly proportional to α_s at leading order. The residual dependency of Eq. (10) upon the strong coupling is thus very well contained. It may be expressed as follows, where A and B are factors that depend only weakly upon α_s , and which may be set by normalization against a pair of known solutions.

$$\begin{aligned} M_{32} &= A \times \exp \left\{ \frac{-2\pi}{\alpha_s (b_2 - b_3)} \right\} \\ \frac{1}{\alpha_5} &= B + \frac{b_2}{\alpha_s (b_2 - b_3)} \end{aligned} \quad (12)$$

The linearized variations may be read off immediately.

$$\begin{aligned} \delta(M_{32}) &= M_{32} \times \frac{-2\pi}{(b_2 - b_3)} \times \delta \left(\frac{1}{\alpha_s} \right) \\ \delta \left(\frac{1}{\alpha_5} \right) &= \frac{b_2}{(b_2 - b_3)} \times \delta \left(\frac{1}{\alpha_s} \right) \end{aligned} \quad (13)$$

Intuitively, lifts in $1/\alpha_s$ are transmitted rather directly to $1/\alpha_5$, modulo some contribution from the differential slopes of the $1/\alpha_2$ and $1/\alpha_3$ logarithmic running. A larger value of α_s will generally require a greater mass scale difference to be spanned in the running, in order to close the wider initial separation between α_2 and α_3 ; as expected, this is manifest in an upward shift of the partial unification point M_{32} . Using the quoted quartic and quadratic scaling of the proton lifetime τ_p with each of M_{32} and $1/\alpha_5$, the linearized variation of this composite factor may also be deduced.

$$\delta(\tau_p) = \tau_p \times \frac{2b_2 \alpha_5 - 8\pi}{(b_2 - b_3)} \times \delta \left(\frac{1}{\alpha_s} \right) \quad (14)$$

The specific β -function coefficients employed in this study are those corresponding to the \mathcal{F} - $SU(5)$ model, as modified by the flippon vector-like multiplet field content.

$$b_2 = 4 \quad ; \quad b_3 = 0 \quad (15)$$

The dominant effect in Eq. (14) is thus the negative term associated with variation of the M_{32} mass scale. In Figure (2), we plot the variation of the proton lifetime with α_s given by the approximation of Eq. Set (12), *i.e.* without feedback into the mass thresholds or second loop tabulation. The absolute numerical scaling of the proton

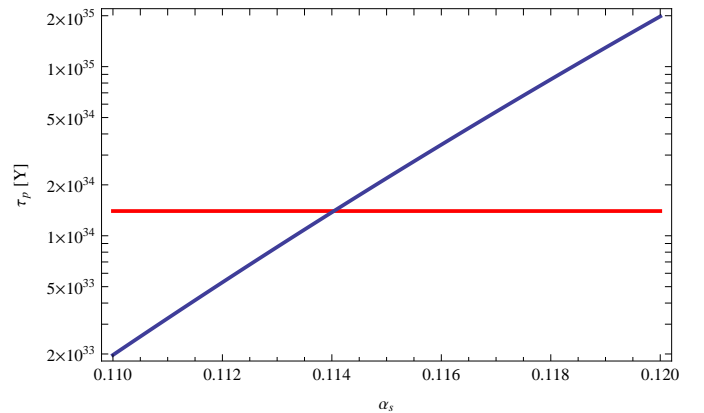


FIG. 2: The variation of the proton lifetime with α_s is depicted by the blue curve. The red horizontal line establishes a lower bound of 1.4×10^{34} Y on the partial proton lifetime in the $e^+ \pi^0$ channel.

lifetime is set relative to the following parameterization which is specialized to flipped $SU(5)$ [35, 56, 57].

$$\tau_p = 3.8 \times \left(\frac{M_{32}}{10^{16} [\text{GeV}]} \right)^4 \times \left(\frac{0.0412}{\alpha_5} \right)^2 \times 10^{35} [\text{Y}] \quad (16)$$

The normalization used is relative to a complete calculation based on a phenomenologically favored benchmark spectrum with a gaugino mass of $M_{1/2} = 708$ GeV, using $\alpha_s = 0.1172$. The results compare very favorably

with a second complete calculation at the same point in the model space, explicitly substituting $\alpha_s = 0.1145$ at all intermediate steps, corresponding to the spectrum provided in Table (1). The isolated values of M_{32} and α_5 track to better than one percent throughout this rather wide gap in α_s , and the proton lifetime agrees within about two percent. The observed variation in the lifetime moreover agrees to a level of about 10 percent across this span with the closed form linearized formula in Eq. (14), given mean values for the referenced τ_p and α_5 baselines. The tracking of $\delta(1/\alpha_5)$ from Eq. Set (13) is less satisfactory, although its influence is subordinate to that of $\delta(M_{32})$, for which the linearized formula is an excellent approximation.

We establish by this parameterization a lower bound of $\alpha_s \geq 0.1140$ for the strong coupling in the context of our favored benchmark $\mathcal{F}\text{-}SU(5)$ model, beneath which the proton lifetime impinges upon the existing limit of 1.4×10^{34} years in the $e^+\pi^0$ decay channel. It should be emphasized that rather fast [35] proton decay is a generic feature of the $\mathcal{F}\text{-}SU(5)$ model, due to intrinsic strengthening of the relevant coupling (α_5) and intrinsic lowering of the relevant mass scale (M_{32}). The situation becomes all the more acute in the context of a mildly depressed input for the strong coupling α_s , as favored both by certain recent direct observations [48] and by the predictions of the the No-Scale $\mathcal{F}\text{-}SU(5)$ model relative to the heaviness of the Higgs mass. The resulting proton lifetime prediction descends in this case to near a point of contact with the current experimental reach, taking a central value of 1.7×10^{34} years for $\alpha_s = 0.1145$. Unaccounted heavy threshold effects near the GUT scale will tend to lengthen rather than shorten this lifetime, by a factor perhaps as large as four [35].

VII. 130 GEV GAMMA-RAY LINE

The observation of a 130 GeV gamma-ray line [58] emanating from our galactic center by the FERMI-LAT Space Telescope has initiated investigations into whether such a monochromatic line could be attributed to dark matter annihilations, an argument amplified by the lack of any known astrophysical source capable of producing a tantamount signature. The lightest neutralino mass at the minimum of the χ^2 fit to the ATLAS multijet and light stop excesses is $m_\chi = 145$ GeV, clearly highlighted as the benchmark model in Figure (1). Conjecturing the observed photon line originates from neutralino annihilations into a Z -boson and gamma-ray via $\tilde{\chi}\tilde{\chi} \rightarrow Z\gamma$, we can compute the kinematics for a non-relativistic lightest neutralino $\tilde{\chi}_1^0$ as

$$E_\gamma = M_\chi - \frac{M_Z^2}{4M_\chi}, \quad (17)$$

which gives

$$M_\chi = \frac{E_\gamma}{2} \left[1 + \sqrt{1 + \left(\frac{M_Z}{E_\gamma} \right)^2} \right] \quad (18)$$

Using $E_\gamma \simeq 130$ GeV and $M_Z = 91.187$ GeV, we arrive at

$$M_\chi = 144.4 \text{ GeV},$$

which is consistent with the No-Scale $\mathcal{F}\text{-}SU(5)$ lightest neutralino mass of $M_\chi = 145$ GeV at the best fit to the multijet and light stop excesses at the LHC and generates an $m_h \simeq 125.5$ GeV Higgs boson mass.

The $\mathcal{F}\text{-}SU(5)$ lightest neutralino composition is greater than 99% bino, therefore, we cannot neglect the quite small $\tilde{\chi}\tilde{\chi} \rightarrow Z\gamma$ annihilation cross-section $\langle\sigma v\rangle \sim 10^{-30}$ cm³/s, about three orders of magnitude less than the FERMI-LAT telescope observations. On the other hand, a recent analysis [59] of extra-galactic clusters uncovering synonymous 130 GeV gamma-ray lines has determined an appropriate subhalo boost factor in this context of ~ 1000 relative to the galactic center. We do not consider it implausible that an overall unaccounted boost factor of similar magnitude might reconcile this apparent discrepancy of scale. For now, we are content to simply make note of the interesting correlation that exists between 145 GeV $\mathcal{F}\text{-}SU(5)$ neutralino annihilations and the unexplained 130 GeV gamma-ray line observed by the FERMI-LAT space telescope, irrespective of the absolute $\langle\sigma v\rangle$ cross-section magnitude.

VIII. IMMINENT TESTABILITY

The full set of correlations enumerated here are imminent testable. We itemized in Ref. [23] the projected signal significance of the ATLAS multijet SUSY searches for the already accumulated 8 TeV collision data, suggesting that indeed an $\mathcal{F}\text{-}SU(5)$ framework presently under probe by the LHC should show evidence in the first 6 fb⁻¹ of 8 TeV data amassed in 2012 of signal significances approaching and possibly breaching the gold standard of $S/\sqrt{B+1} \geq 5$. This tantalizing production of multijet events beyond the Standard Model background has increased proportionally from 1 fb⁻¹ to 5 fb⁻¹ in the 7 TeV collision data, diminishing the prospect of random background fluctuations as the origin of the excess events [23].

We have detailed here the interconnectivity implied by the recently measured $m_h \simeq 125.5 \pm 0.5$ GeV Higgs boson mass and the ATLAS SUSY searches at the LHC. We now append to the markers of imminent testability for No-Scale $\mathcal{F}\text{-}SU(5)$ i) a readily observable proton lifetime, perhaps as low as 1.7×10^{34} years, which is certainly within the reach of the next generation Hyper-Kamiokande and DUSEL experiments [60–62] and already near the grasp of Super-Kamiokande [55], and ii)

a 145 GeV WIMP mass potentially associable with the 130 GeV gamma-ray line observations by the FERMI-LAT Space Telescope.

Whereas the landscape of supersymmetric models is replete with predictions requiring years more of massive data observations and significantly higher LHC beam energies, we emphasize here to the contrary that all the No-Scale \mathcal{F} - $SU(5)$ predictions detailed in this work dwell just on the cusp of an experimentally significant discovery. The mutual correspondence exhibited across multiple independent channels of empirical evidence with the No-Scale \mathcal{F} - $SU(5)$ high-energy framework heightens the potentiality that the coming year could present a complementary discovery to the Higgs boson of equal or even more profound consequence.

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