

Relaxing the upper bound on the mass of the lightest supersymmetric Higgs bosonAndreas Birkedal,¹ Z. Chacko,² and Yasunori Nomura^{3,4}¹*Department of Physics, University of Florida, Gainesville, Florida 32611, USA*²*Department of Physics, University of Arizona, Tucson, Arizona 85721, USA*³*Department of Physics, University of California, Berkeley, California 94720, USA*⁴*Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA*

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We present a class of supersymmetric models in which the lightest Higgs boson mass can be as large as a few hundred GeV (200 ~ 300 GeV) while the successful minimal supersymmetric standard model prediction for gauge coupling unification is preserved. The theories are formulated on a 5D warped space truncated by two branes, and a part of the Higgs sector is localized on the infrared brane. The structure of the Higgs sector in the four-dimensional effective theory below the Kaluza-Klein mass scale is essentially that of the next-to-minimal supersymmetric standard model, or related theories. However, large values of the next-to-minimal supersymmetric standard model couplings at the weak scale are now possible as these couplings are required to be perturbative only up to the infrared cutoff scale, which in general can be much lower than the unification scale. This allows the possibility of generating a large quartic coupling in the Higgs potential thereby significantly raising the Higgs boson mass bound. We present two particularly simple models. In the first model, the quark and lepton fields are localized on the ultraviolet brane, where the grand unified symmetry is broken. In the second model, the quark and lepton fields are localized on the infrared brane, and the unified symmetry is broken both on the ultraviolet and infrared branes. Our theories potentially allow the possibility of a significant reduction in the fine-tuning needed for correct electroweak symmetry breaking, although this is somewhat model dependent.

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I. INTRODUCTION

Postulating supersymmetry as a solution to the hierarchy problem results in significant constraints on the mass of the lightest Higgs boson. For instance, in the minimal supersymmetric standard model (MSSM), the theoretical upper limit on the lightest neutral Higgs boson mass is known to be around 130 GeV, assuming all superpartners have masses below a TeV [1,2]. This is also, roughly, the bound that arises in the minimal extension of the MSSM, so called the next-to-minimal supersymmetric standard model (NMSSM) [3], assuming that the NMSSM singlet couplings remain perturbative up to the scale of gauge coupling unification, Λ_{GUT} [4]. This upper limit is complemented on the lower side by experimental data constraining the lightest Higgs boson mass to be above 114 GeV for many interesting regions of parameter space [5]. Thus, it becomes important to understand the theoretical implications for supersymmetry if the lightest neutral Higgs boson is discovered to have a mass significantly above 130 GeV.

In this paper we present a class of supersymmetric models in which the lightest Higgs boson mass can be as large as about a few hundred GeV (200 ~ 300 GeV) while preserving the successful MSSM prediction for gauge coupling unification. Recall that in the NMSSM the existing tight upper limit for the lightest Higgs mass arises from the requirement that the couplings for the Higgs fields are all perturbative up to the unification scale Λ_{GUT} . However, this requirement does not necessarily have to be imposed if, for example, the Higgs fields are composite (or mixtures

of elementary and composite fields) arising at low energies. In this case the couplings of the singlet field in the NMSSM, for example, can become nonperturbative at a scale much below Λ_{GUT} , significantly weakening the upper bound of the lightest Higgs boson mass. The crucial point is that the low-energy nonperturbative dynamics in the Higgs sector does not necessarily mean that the entire theory enters into a nonperturbative regime below Λ_{GUT} . In fact, it is perfectly possible that the MSSM quark, lepton and gauge sectors stay perturbative up to the scale Λ_{GUT} , allowing a perturbative treatment for gauge coupling evolution, even if the Higgs sector becomes strongly interacting at low energies. Then, as long as the strong dynamics in the Higgs sector respect an approximate global SU(5) symmetry at energies above the TeV scale, the predicted values for the low-energy gauge couplings are the same as those in the MSSM.

How do we explicitly realize the scenario described above? An attractive way of dealing with strong dynamics is to consider higher dimensional theories. Through the Anti-de Sitter/conformal-field-theory (AdS/CFT) duality [6], as applied to the truncated AdS space [7], the above scenario is related to supersymmetric theories in five-dimensional warped space bounded by two branes [8]. The strongly interacting Higgs sector, then, corresponds to the Higgs (and singlet) fields localized to the infrared brane, or the TeV brane, while the perturbative sector corresponds to quarks and leptons (and a part of the Higgs) localized towards the ultraviolet brane, or the Planck brane, with the standard model gauge fields propa-

gating in the bulk. This setup is very attractive because supersymmetry can be broken on the infrared brane, allowing us to naturally understand the hierarchically small supersymmetry-breaking scale through the AdS warp factor [9–14] (see also [15–18]). As shown in [10], this class of theories leaves many of the most attractive features of conventional unification intact; in particular, the successful MSSM prediction for gauge coupling unification is preserved, provided that the 5D bulk possesses an SU(5) gauge symmetry which is broken at the Planck brane and that matter and two Higgs doublets are localized towards the Planck brane or have conformally-flat wave functions. (The successful prediction was anticipated earlier in [19], and techniques for calculating gauge coupling evolution in warped space were developed in [20–25].) These theories also have several nice features if the SU(5) breaking at the Planck brane is caused by boundary conditions. The models we present in this paper preserve the aforementioned attractive features, including the MSSM prediction for gauge coupling unification. Alternative approaches to raising the upper bound on the mass of the lightest supersymmetric Higgs boson have been proposed recently in [26–29]. Earlier work on raising the Higgs mass bound can be found, for example, in Refs. [30,31]. We will comment on the relation of some of these papers to our work in later sections.

A large value for the physical Higgs boson mass has the virtue that it potentially reduces the fine-tuning needed to break the electroweak symmetry at the correct scale in supersymmetric models. It is known that in the MSSM we need a relatively large top-squark mass of $m_{\tilde{t}} \gtrsim 500$ GeV in order to obtain an experimentally allowed physical Higgs boson mass of $m_{\text{Higgs}} \gtrsim 114$ GeV in generic parameter regions. This large top-squark mass then leads to a large soft Higgs mass-squared parameter through radiative corrections given by $m_h^2 \simeq -(3y_t^2/4\pi^2)m_{\tilde{t}}^2 \times \ln(M_{\text{mess}}/m_{\tilde{t}})$, where M_{mess} is the scale at which supersymmetry breaking is mediated to the MSSM sector. This generically leads to fine-tuning of electroweak symmetry breaking as m_h^2 is typically larger than m_{Higgs}^2 by an order of magnitude or larger, especially when $\ln(M_{\text{mess}}/m_{\tilde{t}})$ is large. Because m_{Higgs} can be as large as $200 \sim 300$ GeV in our theory, this tuning could potentially be reduced by a large amount. Moreover, in warped supersymmetric models $\ln(M_{\text{mess}}/m_{\tilde{t}})$ is generically small (see [12]),¹ so that our theory allows relatively large superpartner masses for a given value of m_h^2 , which can be as large as $m_h^2 \simeq m_{\text{Higgs}}^2/2 \simeq (150 \sim 200 \text{ GeV})^2$ without any fine-tuning. In some of the explicit realizations of our theory, the virtue of these properties is somewhat reduced by the fact that there can be large tree-level Higgs soft masses, but we believe it

is significant that we can construct simple models accommodating these features.

The organization of the paper is as follows: In the next section we analyze the NMSSM Higgs sector with the cutoff lowered to $\Lambda \ll \Lambda_{\text{GUT}}$ and see that it can push up the bound on the lightest Higgs boson mass to be as large as about 300 GeV. In Sec. III we construct a theory allowing a lowered cutoff for the Higgs sector while preserving the successful MSSM prediction for gauge coupling unification. The theory is formulated in a 5D supersymmetric warped space, and we also briefly discuss phenomenological consequences of the theory. We can obtain the mass of the lightest Higgs boson as large as about 200 GeV in this theory. In Sec. IV we present an alternative model possessing similar properties but having a different configuration of fields in the warped extra dimension. This model allows the lightest Higgs boson mass as large as 300 GeV, although the model requires an imposition of additional symmetries to be fully realistic. Conclusions are given in Sec. V. Some preliminary results of this paper were presented by one of the authors in [35].

II. LOWERING THE CUTOFF OF THE NMSSM HIGGS SECTOR

The tree-level Higgs boson mass bound in the MSSM is limited by the Z -boson mass m_Z . However, the one-loop corrections from top quarks and squarks are sizable, leading to the upper limit of about 130 GeV. This situation is ameliorated only slightly by going to the NMSSM, which includes an extra gauge singlet S that couples to the Higgs fields through the superpotential:

$$W_{\text{NMSSM}} = \lambda S H_u H_d - \frac{\kappa}{3} S^3 + \text{Yukawa couplings.} \quad (1)$$

Including the singlet couplings in the one-loop Higgs mass bound results in

$$m_{h,1\text{-loop}}^2 \leq m_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta + \frac{3}{4\pi^2} y_t^4 v^2 \sin^4 \beta \ln\left(\frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2}\right), \quad (2)$$

where $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$ are the masses of the two top squarks, m_t is the top-quark mass, y_t is the top Yukawa coupling, $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$, and $v \equiv (\langle H_u \rangle^2 + \langle H_d \rangle^2)^{1/2}$. Here we have set the mixing between the left- and right-handed top squarks to be zero for simplicity and neglected a one-loop correction arising from the coupling λ . At first sight, Eq. (2) appears to be a significant relaxation of the upper bound on the Higgs boson mass. However, the value of λ at the weak scale is generally suppressed if we require λ and κ to remain perturbative up to the scale of gauge coupling unification, $\Lambda_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. Once this requirement is imposed, one finds that the inclusion of the singlet field and its coupling to the Higgs fields increases the upper

¹For alternative ideas to reduce the fine-tuning, see, e.g., [32–34].

limit on the lightest Higgs mass by only about 10 GeV compared with the MSSM [36].

Now we ask what happens if we do not require perturbativity of the couplings in Eq. (1) up to the unification scale. One can see from the renormalization group equations (RGEs) for λ and κ that the value of λ always decreases when it is run down from a high scale:

$$\frac{d\lambda}{d\ln\mu} = \frac{\lambda}{16\pi^2} (4\lambda^2 + 2\kappa^2 + 3y_t^2 + 3y_b^2 + y_\tau^2 - 3g^2 - g'^2), \quad (3)$$

$$\frac{d\kappa}{d\ln\mu} = \frac{6\kappa}{16\pi^2} (\lambda^2 + \kappa^2), \quad (4)$$

where y_b and y_τ are the bottom and tau Yukawa couplings, and g and g' are the $SU(2)_L$ and $U(1)_Y$ gauge couplings. For instance, if we assume that λ becomes nonperturbative at a scale Λ above the weak scale but below Λ_{GUT} , i.e., $\lambda(\Lambda) = 2\pi$, then a lower value of Λ results in a higher value of λ when it is run down to the weak scale. This, therefore, results in a higher upper limit on the mass of the lightest Higgs boson. We have illustrated this in Fig. 1, where we have taken $\kappa(\Lambda) = 0$ and $\lambda(\Lambda) = 2\pi$, which maximizes the Higgs boson mass, and chosen the value of $\tan\beta$ such that the largest Higgs mass is obtained for each value of Λ . A similar analysis was performed earlier in [37]. The improvement gained by lowering the scale at which the Higgs-singlet sector becomes nonperturbative is clear. A scale of $\Lambda = 10^4$ GeV, for example, results in an upper Higgs mass bound of approximately 340 GeV, and lower values of Λ give even larger Higgs boson masses. Although there are many uncertainties in the Higgs mass values obtained in this way, for example, those arising from effects of nonzero values for κ , one-loop effects involving λ , higher-order effects and so on, we expect that we can

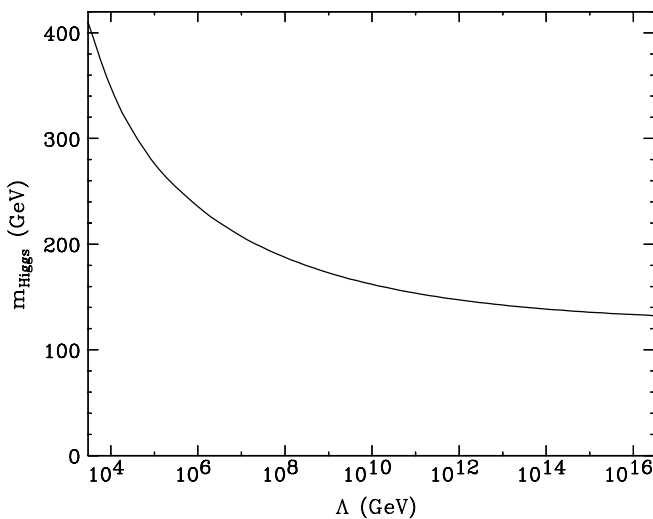


FIG. 1. Theoretical upper limits on the lightest Higgs boson mass as a function of the cutoff scale Λ of the NMSSM Higgs sector.

still obtain the lightest Higgs boson mass as large as about 300 GeV, especially for smaller values of Λ . Note that this result applies to more general superpotentials of the form

$$W_{\text{NMSSM}} = \lambda SH_u H_d + f(S) + \text{Yukawa couplings}, \quad (5)$$

where $f(S)$ is a general function of S : $f(S) = -\Lambda_S^2 S - (M_S/2)S^2 - (\kappa/3)S^3 + \dots$. Here Λ_S and M_S are mass parameters of order the weak scale. The NMSSM is clearly a special case of this.

In deriving the bound on the Higgs boson mass in Fig. 1, we required the top Yukawa coupling to be perturbative ($y_t \lesssim \pi$) only up to the scale Λ and did not require this coupling to be perturbative up to the unification scale of $\Lambda_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. However, we have required all the other sectors of the theory, in particular, the gauge couplings, to remain perturbative up to the unification scale. Our next task is to construct explicit models that realize this type of behavior while preserving the successful MSSM prediction for gauge coupling unification. In the next section, we present a model that shares certain features with the low cutoff NMSSM discussed here, although the resulting upper limits on the Higgs boson mass are tighter than the naive values obtained here. A model that actually realizes the upper limits as high as those given in Fig. 1 will be constructed in Sec. IV.

III. SUPERSYMMETRIC THEORY WITH A HEAVY HIGGS BOSON

A. Basic scheme

In this section we construct a supersymmetric model allowing a relatively heavy (200 GeV or so) Higgs boson while preserving the successful MSSM gauge coupling prediction. We have just seen that a simple way around the tight upper bound on the lightest Higgs boson mass is to allow the Higgs sector to become nonperturbative at a scale below the gauge unification scale Λ_{GUT} . To preserve the perturbative prediction for gauge coupling unification, this must be done in a way such that the sector relevant for the gauge coupling prediction remains perturbative up to Λ_{GUT} . As outlined in the introduction, this can be done by formulating the theory in 5D warped space truncated by two branes.

In a warped extra dimension, the effective cutoff scale of the theory changes with position in the fifth coordinate. We denote this extra dimension by y , where $0 \leq y \leq \pi R$. This can be thought of as arising from compactification on the orbifold S^1/Z_2 . Two branes exist in this setup: an infrared brane at $y = \pi R$ and an ultraviolet brane at $y = 0$. The brane tension causes a warping of the extra dimension into a slice of AdS space. The AdS space describing the extra dimension is defined by the metric:

$$ds^2 = e^{-2k|y|} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2, \quad (6)$$

where k denotes the curvature of the AdS space. The four-

dimensional Planck scale M_4 is then related to the 5D Planck scale M_5 by $M_4^2 \approx M_5^3/2k$ (for $\pi kR \gg 1$). Here we take $k \sim M_5 \sim M_4$, but with k a factor of a few smaller than M_5/π so that the 5D theory is under theoretical control. Identifying M_5 with the cutoff of the 5D theory, the effective y -dependent cutoff, i.e., the cutoff scale measured in terms of the 4D metric $\eta_{\mu\nu}$, is given by

$$\Lambda_{\text{cutoff}} = M_5 e^{-k|y|}. \quad (7)$$

In particular, the effective cutoff on the IR brane is just $\Lambda_{\text{IR}} = M_5 e^{-\pi kR}$, while the cutoff on the UV brane is just the 5D Planck scale, $\Lambda_{\text{UV}} = M_5$. The characteristic scale for the Kaluza-Klein (KK) excitations, which we call the compactification scale, is given by $M_c = \pi k e^{-\pi kR}$. Since we assume $M_5 > \pi k$, the IR cutoff scale is larger than the masses of the first few KK states: $\Lambda_{\text{IR}} > M_c$.

Since the warped extra dimension gives a different effective cutoff at each point in y , fields located in different positions in the extra dimension see different cutoff scales. Since we need our Higgs sector to have a low cutoff scale, we place a part of the Higgs sector, consisting of three chiral superfields S , $H_u^{(b)}$, and $H_d^{(b)}$ and the superpotential of Eq. (5) with H_u and H_d replaced by $H_u^{(b)}$ and $H_d^{(b)}$, on the IR brane (the reason for the superscripts on the Higgs fields will become clear later). The cutoff for the Higgs sector is then given by Λ_{IR} , which is thus identified with the Λ of the previous section. The matter fields are localized on the UV brane so that the cutoff for these fields is around the 4D Planck scale. This suppresses potentially dangerous proton decay and flavor-changing processes. An important point is that, although the Higgs sector becomes nonperturbative at a low-energy scale of Λ_{IR} , it does not affect the physics of the rest of the model. In particular, the evolution of the gauge couplings is not affected by this nonperturbativity of the Higgs sector. This is because physics above the scale M_c measured by a Planck brane observer is not affected by physics on the IR brane [21,38]. For example, a momentum mode p will only feel physics on the IR brane with a strength proportional to $e^{-\pi p/M_c}$. Since $\Lambda_{\text{IR}} > M_c$, any nonperturbative physics at Λ_{IR} in the Higgs sector will have decoupled from physics described by momenta $p > M_c$ on the Planck brane. This is the key feature of AdS space that allows IR physics to become nonperturbative without affecting physics on the UV brane above the scale M_c . Any running of gauge couplings above M_c will not feel any (possibly nonperturbative) physics on the IR brane.

There remain two issues for model-building. How can we transmit the electroweak symmetry breaking caused by the vacuum expectation values (VEVs) of $H_u^{(b)}$ and $H_d^{(b)}$ on the TeV brane to the quarks and leptons localized on the Planck brane? And how can we maintain the MSSM prediction for gauge coupling unification? Both of these issues are simultaneously resolved if we introduce two additional Higgs doublets $H_u^{(B)}$ and $H_d^{(B)}$ in the bulk.

These fields can interact both with the Planck and TeV branes and can transmit electroweak symmetry breaking. Further, the addition of these fields is sufficient to preserve the MSSM prediction for gauge coupling unification, as we will see later.

B. Model

We now construct our model. While most of our construction applies to general values of Λ_{IR} , we concentrate on the case where $\Lambda_{\text{IR}} \sim \text{TeV}$, i.e., $kR \sim 10$, in what follows because it gives the largest upper bound on the Higgs boson mass and allows a simple implementation of supersymmetry breaking. We take the gauge group in the bulk to be $\text{SU}(5)$, which is broken by boundary conditions at the Planck brane [10]. Specifically, the 5D gauge supermultiplet, which can be decomposed into a 4D $N = 1$ vector superfield V and a 4D $N = 1$ chiral superfield Σ , obeys the following set of boundary conditions:

$$\begin{aligned} \begin{pmatrix} V \\ \Sigma \end{pmatrix}(x^\mu, -y) &= \begin{pmatrix} PVP^{-1} \\ -P\Sigma P^{-1} \end{pmatrix}(x^\mu, y), \\ \begin{pmatrix} V \\ \Sigma \end{pmatrix}(x^\mu, -y') &= \begin{pmatrix} P'VP'^{-1} \\ -P'\Sigma P'^{-1} \end{pmatrix}(x^\mu, y'), \end{aligned} \quad (8)$$

where $y' = y - \pi R$. Here, V and Σ are both in the adjoint of $\text{SU}(5)$, and P and P' are 5×5 matrices acting on gauge space taken here as $P = \text{diag}(+, +, +, -, -)$ and $P' = \text{diag}(+, +, +, +, +)$. This reduces the gauge symmetry at low energies to be $\text{SU}(3)_C \times \text{SU}(2)_L \times \text{U}(1)_Y$ (321): only the 321 component of V has a zero mode. The characteristic scale for the KK tower is M_c , which is a factor of a few smaller than the IR cutoff scale Λ_{IR} .

We also introduce two bulk hypermultiplets $\{H, H^c\}$ and $\{\bar{H}, \bar{H}^c\}$ in the fundamental representation of $\text{SU}(5)$. Here, we have decomposed a hypermultiplet into two 4D $N = 1$ chiral superfields, where $H(\mathbf{5})$, $H^c(\mathbf{5}^*)$, $\bar{H}(\mathbf{5}^*)$, $\bar{H}^c(\mathbf{5})$ are 4D chiral superfields with the numbers in parentheses representing their transformation properties under $\text{SU}(5)$. The boundary conditions are given by

$$\begin{aligned} \begin{pmatrix} H \\ H^c \end{pmatrix}(x^\mu, -y) &= \begin{pmatrix} -PH \\ PH^c \end{pmatrix}(x^\mu, y), \\ \begin{pmatrix} H \\ H^c \end{pmatrix}(x^\mu, -y') &= \begin{pmatrix} P'H \\ -P'H^c \end{pmatrix}(x^\mu, y'), \end{aligned} \quad (9)$$

for $\{H, H^c\}$, and similarly for $\{\bar{H}, \bar{H}^c\}$. The zero modes then arise only from the $\text{SU}(2)_L$ -doublet components of H and \bar{H} , which we call $H_u^{(B)} \subset H$ and $H_d^{(B)} \subset \bar{H}$. In general, a bulk hypermultiplet $\{\Phi, \Phi^c\}$ can have a mass term in the bulk, which is written as

$$S = \int d^4x \int_0^{\pi R} dy \left[e^{-3k|y|} \int d^2\theta c_\Phi k \Phi \Phi^c + \text{H.c.} \right], \quad (10)$$

in the basis where the kinetic term is given by $S_{\text{kin}} = \int d^4x \int dy [e^{-2|k|y} \int d^4\theta (\Phi^\dagger \Phi + \Phi^c \Phi^{c\dagger}) + \{e^{-3|k|y} \times \int d^2\theta (\Phi^c \partial_y \Phi - \Phi \partial_y \Phi^c)/2 + \text{H.c.}\}]$ [39]. The parameter c_Φ controls the wave function profile of the zero mode— for $c_\Phi > 1/2$ ($< 1/2$) the wave function of a zero mode arising from Φ is localized to the Planck (TeV) brane; for $c_\Phi = 1/2$ it is conformally flat. For H and \bar{H} fields, we choose c_H and $c_{\bar{H}}$ to be $1/2$ (or slightly larger than $1/2$):

$$c_H \simeq c_{\bar{H}} \simeq \frac{1}{2}, \quad (11)$$

so that the zero modes of $H_u^{(B)}$ and $H_d^{(B)}$ have (almost) conformally-flat wave functions. These fields then play a role of transmitting electroweak symmetry breaking in the TeV-brane Higgs sector, which will be introduced later, to the matter sector ($H_u^{(B)}$ and $H_d^{(B)}$ are not completely identified with the Higgs doublets in the MSSM, as we will see below). The matter fields, Q , U , D , L , and E for each generation, are introduced on the Planck brane with the Yukawa couplings to $H_u^{(B)}$ and $H_d^{(B)}$:

$$S = \int d^4x \int_0^{\pi R} dy 2\delta(y) \left[\int d^2\theta (y_u Q U H_u^{(B)} + y_d Q D H_d^{(B)} + y_e L E H_d^{(B)}) + \text{H.c.} \right]. \quad (12)$$

(Alternatively, matter could be introduced in the bulk with the appropriate boundary conditions and wave functions localized towards the Planck brane by having $c \gg 1/2$. All our analyses apply to both the brane and bulk matter cases.) With the above configuration of fields, the prediction for the low-energy gauge couplings are the same as in the MSSM at the leading order and proton decay rates are adequately suppressed [10]. Small neutrino masses are also naturally obtained through the see-saw mechanism by introducing right-handed neutrinos N on the Planck brane or in the bulk [10].

We now introduce the Higgs sector on the TeV brane, which could become strongly coupled at the IR cutoff scale $\Lambda_{\text{IR}} \sim \text{TeV}$. We introduce three chiral superfields $S(\mathbf{1})$, $H'(\mathbf{5})$, and $\bar{H}'(\mathbf{5}^*)$ on the TeV brane with the numbers in parentheses representing the transformation properties under $SU(5)$ [the active gauge group on the TeV brane is $SU(5)$ so that any multiplet on this brane must be in a representation of $SU(5)$]. We introduce the superpotential of the following form on the brane:

$$S = \int d^4x \int_0^{\pi R} dy 2\delta(y - \pi R) \left[e^{-3\pi k R} \int d^2\theta (\lambda S H' \bar{H}') + f(S) + \text{H.c.} \right], \quad (13)$$

where, as explained earlier, $f(S)$ is a function of S with the general form given by $f(S) = -\Lambda_S^2 S - (M_S/2)S^2 -$

$(\kappa/3)S^3 + \dots$. Now, however, Λ_S and M_S are mass parameters of the order of M_5 , or somewhat smaller, in the original 5D metric. Since the sector living on the IR/TeV brane only needs to stay perturbative up to Λ_{IR} , we only need to require $\lambda(\Lambda_{\text{IR}}), \kappa(\Lambda_{\text{IR}}) \lesssim 2\pi$. This gives a large quartic coupling for the doublet components of H' and \bar{H}' , which we call $H_u^{(b)} \subset H'$ and $H_d^{(b)} \subset \bar{H}'$ (these fields consist of parts of the MSSM Higgs doublets, as we will see shortly). Once supersymmetry is broken, $H_u^{(b)}$ and $H_d^{(b)}$ (and S) will obtain VEVs, breaking the electroweak symmetry. An important point is that, due to the properties of AdS space, the introduction of TeV-brane fields and/or superpotentials does not modify the physics at higher energies, including the evolution of the gauge couplings. Therefore, the prediction for gauge coupling unification is still the same as that in the MSSM.

To complete the construction of the model, we have to connect the two sets of Higgs doublets, $H_u^{(B)}, H_d^{(B)}$ and $H_u^{(b)}, H_d^{(b)}$. This can be done by introducing mixing terms between these fields on the TeV brane:

$$S = \int d^4x \int_0^{\pi R} dy 2\delta(y - \pi R) \left[e^{-3\pi k R} \int d^2\theta (M_{ud}^{1/2} H \bar{H}' + M_{du}^{1/2} \bar{H} H') + \text{H.c.} \right], \quad (14)$$

where M_{ud} and M_{du} are parameters having the mass dimension of 1 and of order M_5 . Note that the doublet components of H , \bar{H} , H' , and \bar{H}' are $H_u^{(B)}, H_d^{(B)}, H_u^{(b)}$, and $H_d^{(b)}$, respectively. In the KK-decomposed 4D theory, interactions of Eqs. (13) and (14) give the following supersymmetric masses for the Higgs doublets

$$W = \begin{pmatrix} H_u^{(B)} & H_u^{(b)} \end{pmatrix} \begin{pmatrix} 0 & e^{-\pi k R} \sqrt{\frac{M_{ud}}{\pi R}} \\ e^{-\pi k R} \sqrt{\frac{M_{du}}{\pi R}} & \lambda \langle S \rangle \end{pmatrix} \begin{pmatrix} H_d^{(B)} \\ H_d^{(b)} \end{pmatrix}, \quad (15)$$

where (and below) $H_u^{(B)}$ and $H_d^{(B)}$ stand for the zero modes for these fields, and S is canonically normalized in 4D. Here, we have assumed that M_{ud} and M_{du} are of order M_5 or smaller and taken into account the volume-suppression factor arising from wave functions of $H_u^{(B)}$ and $H_d^{(B)}$. We have also left out the potential coupling of the singlet to the bulk Higgs fields which is more volume suppressed than the other terms and therefore negligible. This gives the desired mixing between $H_u^{(B)}$ and $H_u^{(b)}$, and $H_d^{(B)}$ and $H_d^{(b)}$. After supersymmetry is broken, a certain parameter region of the model leads to VEVs for one set of the Higgs doublets H_u and H_d parameterized by

$$H_u = \cos\theta_u H_u^{(b)} + \sin\theta_u H_u^{(B)}, \quad (16)$$

$$H_d = \cos\theta_d H_d^{(b)} + \sin\theta_d H_d^{(B)}, \quad (17)$$

which we assume to be the case. These fields H_u and H_d , therefore, are the Higgs doublets responsible for electro-weak symmetry breaking. Assuming that S does not have a large supersymmetric mass, they have a large quartic term arising from the superpotential coupling $W = \lambda \cos\theta_u \cos\theta_d S H_u H_d$:

$$V_H = \lambda^2 \cos^2\theta_u \cos^2\theta_d |H_u H_d|^2. \quad (18)$$

Furthermore, these doublets also couple to the quarks and leptons through the Yukawa couplings of Eq. (12):

$$W = y'_u \sin\theta_u Q U H_u + y'_d \sin\theta_d Q D H_d + y'_e \sin\theta_d L E H_d, \quad (19)$$

giving the quark and lepton masses, where $y'_{u,d,e} \equiv y_{u,d,e}/(\pi R)^{1/2}$ are dimensionless coupling parameters. Therefore, for reasonable $O(1)$ values for the mixing angles, interactions in Eqs. (18) and (19) can give the required quark-lepton masses and a large quartic term for the Higgs fields.

What if the IR scale Λ_{IR} is much larger than the TeV scale? In this case M_{ud} and M_{du} [and Λ_S and M_S in the function $f(S)$] must be chosen such that their values measured in terms of the 4D metric are of order TeV, i.e., $e^{-\pi k R} M_{ud} \sim e^{-\pi k R} M_{du} \sim \text{TeV}$. Alternatively, for the NMSSM, one could introduce an additional set of Higgs fields H'' and \bar{H}'' on the TeV brane with the superpotential term $\delta(y - \pi R)\{\hat{M}_{ud}^{1/2} H''(a\bar{H} + b\bar{H}') + \hat{M}_{du}^{1/2} \bar{H}''(cH + dH')\}$, in which case \hat{M}_{ud} and \hat{M}_{du} could be as large as M_5 .

In Fig. 2 we have plotted the upper limit on the physical Higgs boson mass as a function of Λ . In the figure, we have set the mixing angle for the down-type Higgs to be zero, $\theta_d = 0$, for simplicity, and plotted the limits for various different values of the mixing angle for the up-type Higgs θ_u . The bound plotted is maximized by setting $\lambda(\Lambda) = 2\pi$ and $\kappa(\Lambda) = 0$, where $\Lambda \equiv \Lambda_{\text{IR}}$, and by using the model with additional Higgs doublets H'' and \bar{H}'' described in the previous paragraph for higher values of Λ . In the figure, we have applied the condition of $\tan\beta > 2$, because the values of $\tan\beta$ much smaller than 2 require very small $\cos\theta_u$ to obtain a large enough top-quark mass and thus lead to a small physical Higgs boson mass (the maximum Higgs boson mass may actually be obtained at $\tan\beta$ somewhat smaller than 2, but the bound does not change much even in that case). We have also assumed that the mixing of the two sets of Higgs doublets, Eqs. (16) and (17), occurs at the scale Λ . This maximizes a Higgs boson mass obtained for a given value of Λ and θ_u . In a realistic situation, however, the scale at which the Higgs mixing occurs will most likely be a factor of a few smaller than Λ . This could give a reduction of the Higgs boson mass as large as about (30 ~ 50) GeV for $\Lambda \sim \text{TeV}$ depending on details of the physics at $\sim \Lambda$, since the coupling λ is then very large between Λ and the scale of the Higgs mixing, which could give large negative corrections to the low-energy values of λ and the

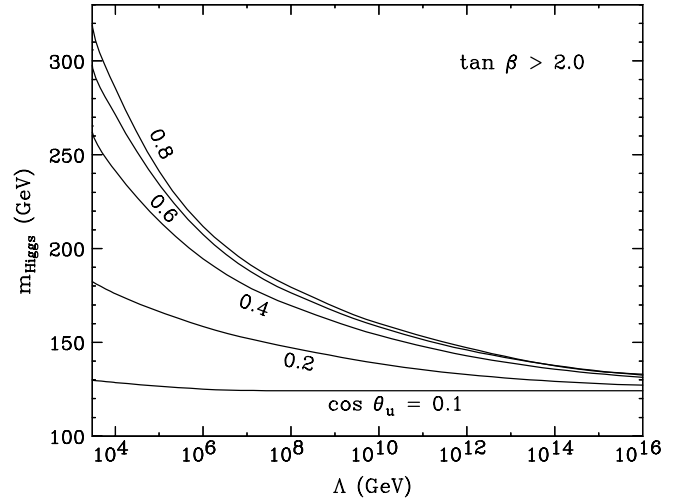


FIG. 2. Upper limits on the lightest Higgs boson mass as a function of the IR cutoff scale Λ for various values of the up-type Higgs mixing angle θ_u : $\cos\theta_u = 0.1, 0.2, 0.4, 0.6,$ and 0.8 .

top Yukawa coupling. The reduction is smaller for larger values of Λ .

The value of $\cos\theta_u$ is bounded from above by requiring that a large enough top-quark mass is obtained through the Yukawa coupling of Eq. (19). For example, if one neglects the effect of the running of y_u (but not the volume-suppression effect), one finds that the upper bound on $\cos\theta_u$ is about 0.6 for $\Lambda \approx 10$ TeV and $\tan\beta \gtrsim 2$. The running effect between Λ and the electroweak scale could somewhat reduce this value, e.g., to $\cos\theta_u \lesssim 0.55$, but we still see from Fig. 2 that our theory allows the Higgs boson mass as high as 200 GeV, even after considering the potential reduction of the Higgs mass coming from physics at the scale Λ . To complete the discussion, we must also consider the running of the top Yukawa coupling y_u above the scale Λ . We will, however, see at the end of this section that the Higgs mass bound obtained here is not much changed by this effect. Because our theory admits a large physical Higgs boson mass, it potentially allows a significant reduction (or an elimination) of the fine-tuning.

Let us here consider an example of the parameter region leading to a large Higgs boson mass. We here concentrate on the case with $\Lambda_{\text{IR}} = O(10 \text{ TeV})$ and the region where the elements of the Higgs mass matrix in Eq. (15) are smaller than M_c . In this parameter region mixings between the light Higgs states and higher KK states are negligible, so that we can neglect the effect of the KK states in analyzing electroweak symmetry breaking (note also that the precision electroweak constraints from the KK states are very weak with matter localized on the Planck brane [40]). For $M'_{du} \lesssim M'_{ud} \approx \lambda\langle S \rangle$, where $M'_{xy} \equiv e^{-\pi k R} \sqrt{M_{xy}/\pi R}$ with $x, y = u, d$, we obtain $\theta_u = O(1)$ and $\theta_d \lesssim 1$. This leads to somewhat suppressed Yukawa couplings for the down-type quarks and charged leptons

through Eq. (19), but not for the up-type quarks (especially the top quark). The tree-level Higgs-quartic couplings are given by Eq. (18) plus the contribution from the $SU(2)_L \times U(1)_Y$ D terms. The masses for two sets of Higgs doublets are of order $M_1 = M'_{ud}M'_{du}/\lambda\langle S \rangle$ and $M_2 = \lambda\langle S \rangle$, where $M_1 \lesssim M_2$ (the mass of the Higgs triplets arising from H' and \bar{H}' is of order M_2). Supersymmetry breaking is caused on the TeV brane by a nonzero VEV of a chiral superfield Z : $F_Z = \langle \partial_{\theta^2} Z \rangle \neq 0$, which gives masses for the gauginos through the operator $\delta(y - \pi R) \int d^2\theta Z \mathcal{W}^\alpha \mathcal{W}_\alpha$ and for the squarks and sleptons through finite loop contributions.² Then, if supersymmetry-breaking masses for the Higgs fields, arising from operators of the form $\delta(y - \pi R) \sum_{i,j} \int d^4\theta Z^\dagger Z (H_i \bar{H}_j + H_i^\dagger H_j + \bar{H}_i^\dagger \bar{H}_j)$ where $H_i = H, H'$ and $\bar{H}_i = \bar{H}, \bar{H}'$ are of order M_1 , correct electroweak symmetry breaking can be induced. Supersymmetry breaking should also trigger a nonzero VEV for S , and we assume that the supersymmetric mass for S —e.g., $\kappa\langle S \rangle$ in the case with $M_S = 0$ —is of order M_1 . Such a VEV, for example, may arise if the soft supersymmetry-breaking Lagrangian (and thus the superpotential) contains a linear term in S , $\mathcal{L}_{\text{soft}} = (C_S \Lambda_S^2 S + \text{H.c.}) + \dots$, in which case the VEV of S is given by $\langle S \rangle \simeq C_S \Lambda_S^2 / m_S^2$ in certain parameter region, where m_S^2 is the soft supersymmetry-breaking mass squared for S . This allows relatively small values of $\langle S \rangle$ even with large values of m_S^2 . The spectrum contains an extra pair of Higgs doublets and a pair of Higgs triplets with the masses of order M_2 , in addition to the states in the NMSSM. In fact, the presence of an extra pair of Higgs fields with the quantum numbers of $\mathbf{5} + \mathbf{5}^*$ under $SU(5)$ is a generic prediction of the model with $\Lambda_{\text{IR}} \sim \text{TeV}$.

Finally, we consider the issue of evolution of the top Yukawa coupling above Λ_{IR} : y_u in Eq. (12) for the third generation. In the present model the evolution of the Yukawa coupling receives additional contribution from the bulk, which could potentially alter the existence and location of a Landau pole for the top Yukawa coupling (and for the other couplings). While we do not make a full analysis of these coupling evolutions, we can make a rough estimate of this effect in the following way. We first rescale the 5D Higgs field $H_u^{(B)}$ as $H_u^{(B)} \rightarrow \sqrt{M} \hat{H}_u$, so that \hat{H}_u has a mass dimension of 1. Then the top Yukawa coupling of Eq. (12) is written as $2\delta(y) \int d^2\theta y_u \sqrt{M} Q U \hat{H}_u + \text{H.c.}$ (note that y_u has mass dimensions of $-1/2$). Suppose that the fields Q , U , and \hat{H}_u have brane-localized kinetic

terms of $2\delta(y) \int d^4\theta (Z_{0,Q} Q^\dagger Q + Z_{0,U} U^\dagger U + Z_{0,H} \hat{H}^\dagger \hat{H})$ at tree level. In the dual 4D picture, this implies that the wave function renormalizations for the fields Q , U , and \hat{H} , defined as $Z_Q(\mu)$, $Z_U(\mu)$, and $Z_H(\mu)$, take the values $Z_{0,Q}$, $Z_{0,U}$, and $Z_{0,H}$ at the scale k :

$$Z_Q(k) = Z_{0,Q}, \quad Z_U(k) = Z_{0,U}, \quad Z_H(k) = Z_{0,H}, \quad (20)$$

where μ is the renormalization scale. When we evolve the RGEs to low energies, Z_Q , Z_U , and Z_H receive quantum corrections (but $y_u \sqrt{M}$ does not receive such corrections due to the nonrenormalization theorem). In particular, $Z_H(\mu = k e^{-\pi k R})$ receives a contribution from the bulk, which we interpret as the running effect: $\delta Z_{H,\text{bulk}} = \pi R M = (M/k) \ln(k/\mu)$. Then, if we simply add the bulk contribution to the MSSM running, we obtain RGEs for the wave functions:

$$\frac{d \ln Z_Q}{d \ln \mu} = -\frac{1}{8\pi^2} \left(\frac{y_u^2 M}{Z_Q Z_U Z_H} - \frac{8}{3} g_3^2 \right), \quad (21)$$

$$\frac{d \ln Z_U}{d \ln \mu} = -\frac{1}{8\pi^2} \left(2 \frac{y_u^2 M}{Z_Q Z_U Z_H} - \frac{8}{3} g_3^2 \right), \quad (22)$$

$$\frac{d \ln Z_H}{d \ln \mu} = -\frac{1}{8\pi^2} \left(3 \frac{y_u^2 M}{Z_Q Z_U Z_H} \right) - \frac{M}{k Z_H}. \quad (23)$$

The low-energy top Yukawa coupling y'_t , which couples Q and U with the zero mode of $H_u^{(B)}$ [the 33 element of y'_u in Eq. (19)], is then obtained as

$$y'_t(\mu)^2 = \frac{y_u^2 M}{Z_Q(\mu) Z_U(\mu) Z_H(\mu)}. \quad (24)$$

Note that $Z_H(\mu)$ [$Z_Q(\mu)$ and $Z_U(\mu)$] is proportional to M [M^0] so that the coupling y'_t does not depend on the spurious parameter M . Using these RGEs, we can obtain the low-energy value of y'_t . Note that the $SU(3)_C$ gauge coupling g_3 obeys the RGE with the bulk contribution added. Assuming that the Planck brane-localized kinetic term at tree level is small, the RGE takes the form $d(1/g_3^2)/d \ln \mu = -(b/8\pi^2)$ with $b = b_{\text{MSSM}} + b_{\text{bulk}} \simeq 1.8$, where b_{bulk} represents the $SU(5)$ -invariant bulk contribution, which makes g_3 nonperturbative at the scale k . For natural sizes for the coefficients in 5D, i.e., $Z_{0,Q} \sim Z_{0,U} \sim 1$, $Z_{0,H} \sim M/M_*$, and $y_u \sim 4\pi/\sqrt{M_*}$, where M_* is the 5D cutoff scale, we obtain $y'_t(\mu \sim \text{TeV}) \lesssim (1.3 \sim 1.4)$ for $\Lambda_{\text{IR}} \sim \text{TeV}$. Although this estimate is somewhat uncertain, we expect that we can obtain a large enough top-quark mass $m_t = y'_t \sin\theta_u \langle H_u \rangle$, where $\langle H_u \rangle = v \sin\beta$, for $\sin\theta_u$ not much smaller than 1, e.g., $\sin\theta_u \gtrsim 0.8$ for $\tan\beta \gtrsim 2$. This in turn gives the upper limit on the lightest Higgs boson mass larger than 200 GeV for small values of Λ_{IR} , since $\cos\theta_u$ as large as 0.6 is allowed (see Fig. 2). Although a more careful analysis may be needed to be

²An alternative possibility is to break supersymmetry on the Planck brane with an intermediate scale VEV for F_Z , so that the 321 gauginos receive weak scale masses on the Planck brane. In this case, the Higgs fields H' and \bar{H}' do not obtain tree-level supersymmetry-breaking masses (and the supersymmetry-breaking masses for $H_u^{(B)}$ and $H_d^{(B)}$ are volume suppressed). The squark and slepton masses, however, are generated at tree-level through the couplings to Z , which must be assumed to be flavor universal.

really conclusive about the issue of the top-quark mass, we expect that a Higgs boson mass as large as 200 GeV can be obtained in this model.³

IV. MODEL WITH MATTER ON THE IR BRANE

In this section we construct an alternative model. This model differs from that of the previous section in the location of fields. In particular, we locate the quark and lepton fields on the IR brane so that there is no issue of nonstandard evolution of the Yukawa couplings (the evolution of the Yukawa couplings is the standard one in 4D below Λ_{IR}). Because the Yukawa couplings are located on the IR brane, they can become nonperturbative at the scale Λ_{IR} , giving the observed top-quark mass quite easily. We again concentrate on the case with $\Lambda_{\text{IR}} \sim \text{TeV}$ below, since it gives the largest bound on the Higgs boson mass and the simplest realization of gauge coupling unification.⁴

The model uses a bulk SU(5) gauge symmetry, but it is now broken at both the Planck and the TeV branes, i.e., the 5D gauge multiplet obeys the boundary conditions of Eq. (8) with $P = P' = \text{diag}(+, +, +, -, -)$ [13]. The bulk Higgs fields are introduced as two hypermultiplets in the fundamental representation of SU(5), $\{H, H^c\}$, and $\{\bar{H}, \bar{H}^c\}$, with the boundary conditions given by Eq. (9) and $P = P' = \text{diag}(+, +, +, -, -)$. These boundary conditions yield zero modes that are not the MSSM states, the SU(5)/[SU(3)_C × SU(2)_L × U(1)_Y] component of Σ , in addition to the MSSM gauge fields, the 321 component of V . However, these unwanted states can be made heavy (with masses of order TeV) through supersymmetry-breaking effects on the TeV brane, and the prediction for the low-energy gauge couplings is still that of the MSSM, as long as the bulk mass parameters for the Higgs fields, c_H

and $c_{\bar{H}}$, satisfy $c_H, c_{\bar{H}} \geq 1/2$ [13]. In this “321-321 model,” the gauge groups effective at the Planck brane and the TeV brane are both SU(3)_C × SU(2)_L × U(1)_Y. Therefore, we can locate the MSSM matter fields Q , U , D , L , and E on the TeV brane without introducing proton decay mediated by the SU(5) gauge bosons. Potentially dangerous tree-level proton decay operators can also be forbidden by imposing a symmetry, say baryon number, on the model.⁵

We now take a closer look at the Higgs sector. For the boundary conditions described above, the Higgs fields do not possess zero modes. However, if the bulk mass parameters for the Higgs fields are much larger than $k/2$, i.e., $c_H, c_{\bar{H}} \gg 1/2$, four doublet states from H , H^c , \bar{H} , and \bar{H}^c become exponentially lighter than M_c , which we assume to be the case (see footnote 2 of [13]). Among these states, the modes arising from H and \bar{H} , which we call H'_u and H'_d , respectively, are (strongly) localized towards the Planck brane, while those arising from H^c and \bar{H}^c , which we call H_d and H_u , respectively, are (strongly) localized towards the TeV brane. Now, we can introduce superpotential interactions between the H_u and H_d fields and the fields located on the TeV brane. Specifically, we introduce Yukawa couplings between the H_u and H_d fields and the quarks and leptons:

$$S = \int d^4x \int_0^{\pi R} dy 2\delta(y - \pi R) \left[e^{-3\pi k R} \int d^2\theta (y'_u Q U H_u + y'_d Q D H_d + y'_e L E H_d) + \text{H.c.} \right]. \quad (25)$$

In addition, we introduce a singlet chiral superfield S on the TeV brane together with the superpotential couplings

$$S = \int d^4x \int_0^{\pi R} dy 2\delta(y - \pi R) \left[e^{-3\pi k R} \int d^2\theta (\lambda' S H_u H_d + f(S)) + \text{H.c.} \right], \quad (26)$$

where $f(S)$ is as before. Since the light H_u and H_d fields are strongly localized towards the TeV brane, the couplings in Eqs. (25), (26) do not receive a volume-suppression factor when reduced to the low-energy 4D theory. This allows large couplings for the low-energy superpotential

³It is possible to construct models in which the running of the top Yukawa coupling is the standard 4D running. An example of such models is the following: We have TeV-brane Higgs fields S , H' , and \bar{H}' together with the superpotential of Eq. (13). We further introduce a Higgs field H''_u on the Planck brane and a Higgs hypermultiplet $\{\bar{H}, \bar{H}^c\}$ in the bulk, whose zero mode is denoted as $H_d^{(B)}$. These fields are coupled to matter fields localized on the Planck brane as Eq. (12), but with $H_u^{(B)}$ replaced by H''_u . Then, introducing the Higgs mixing terms of the form $\delta(y) \int d^2\theta H''_u H_d^{(B)} + \delta(y - \pi R) \int d^2\theta H_u^{(B)} H_d^{(B)}$, we obtain the top Yukawa and Higgs-quartic couplings at low energies. To get the 4D running for the top Yukawa coupling, however, we must arrange the coefficients of the two Higgs mixing terms such that they both have sizes around Λ_{IR} when measured in terms of the 4D metric $\eta_{\mu\nu}$. Simple generalizations of this idea also lead to realistic down-type Yukawa couplings, while preserving the prediction for gauge coupling unification.

⁴The Higgs boson mass bound in warped supersymmetric theories with matter fields on the TeV brane was also considered in [31]. The model discussed there, however, does not accommodate the MSSM prediction for gauge coupling unification, and the bound on the Higgs mass is relaxed only for low values of Λ of order TeV, as it uses operators suppressed by powers of Λ .

⁵We comment here that it is straightforward to use the 321-321 model to construct a theory of the type discussed in Sec. III, i.e., the model with the quarks and leptons on the Planck brane. The simplest of such models has $\Lambda_{\text{IR}} \sim \text{TeV}$, with the boundary conditions for the gauge multiplet given by Eq. (8) with $P = P' = \text{diag}(+, +, +, -, -)$ and those for the bulk Higgs hypermultiplets given by Eq. (9) with $P = -P' = \text{diag}(+, +, +, -, -)$. The TeV-brane Higgs sector consists of S , $H_u^{(b)}$, and $H_d^{(b)}$ with the superpotential interactions of the form $\delta(y - \pi R) \int d^2\theta \{\lambda S H_u^{(b)} H_d^{(b)} + f(S)\}$. The rest of the constructions of the model is analogous to those of Sec. III.

$$W = \lambda S H_u H_d + f(S) + y_u Q U H_u + y_d Q D H_d + y_e L E H_d, \quad (27)$$

where $y_{u,d,e} = y'_{u,d,e} \sqrt{k}$ and $\lambda = \lambda' \sqrt{k}$. Since λ in Eq. (27) can be large of order π at low energies of $E \simeq M_c$, a large physical Higgs boson mass can be obtained. The upper limits on the Higgs boson mass are essentially given by Fig. 1 because there is no issue in obtaining a large enough top-quark mass. Because of the bound on the masses of the KK states [40], we expect that $\Lambda_{\text{IR}} \gtrsim 100$ TeV in the present model. However, this still allows a Higgs boson mass as large as about 300 GeV.

There are several issues in the present model. First, small neutrino masses can be generated by introducing right-handed neutrino hypermultiplets $\{N, N^c\}$ in the bulk with the $(+, +)$ and $(-, -)$ boundary conditions for N and N^c , respectively, and coupling them to the lepton doublets on the TeV brane through the operators of the form $\delta(y - \pi R) \int d^2\theta L N H_u$. Then, if the bulk masses for right-handed neutrinos are large, i.e., $c_N \gg 1/2$, we can naturally obtain small (Dirac) neutrino masses [41]. Potentially large neutrino masses from the TeV-brane operators $\delta(y - \pi R) \int d^2\theta (L H_u)^2$ can be forbidden if we impose lepton number with charges $L(1), N(-1)$ on the model. Second, the H'_u and H'_d states, which are localized towards the Planck brane, must obtain a mass of order Λ_{IR} through the superpotential term $\delta(y) \int d^2\theta H'_u H'_d$ to evade phenomenological constraints and to preserve the MSSM prediction for gauge coupling unification. The required mass term can naturally be generated by shining the scale from the TeV to Planck brane through a bulk singlet field with $c \simeq 1/2$, as discussed in [10]. Alternatively, for $\Lambda_{\text{IR}} \sim$ TeV, the mass of the H'_u and H'_d states of order Λ_{IR} can be generated through supersymmetry breaking by introducing a singlet field X on the Planck brane with the superpotential interaction of the form $\delta(y) \times \int d^2\theta (X H'_u H'_d + X^3)$, as in the NMSSM. Finally, supersymmetry breaking on the TeV brane will give masses at tree level not only to the gauginos but also to the squarks and sleptons through the operators of the form $\delta(y - \pi R) \int d^4\theta Z^\dagger Z (Q^\dagger Q + U^\dagger U + D^\dagger D + L^\dagger L + E^\dagger E)$. This in general leads to the supersymmetric flavor problem, which we somehow have to avoid. This may be accomplished, for example, by imposing a flavor symmetry such as $U(2)_F$ acting on the first two generations to these interactions. Such a symmetry should be broken to generate realistic quark and lepton mass matrices, which in turn gives a small nonuniversality in the squark and slepton mass matrices. Although it is not obvious that the supersymmetric flavor problem is naturally solved along this line, here we do not attempt to make detailed studies on this issue and simply assume that it can be done in a phenomenologically successful manner (if not, we have to make an assumption for the interactions between the Z field and the quarks and leptons, which is not explained in

our effective field theory). Since supersymmetry-breaking masses for the gauginos, sfermions, and Higgs fields can naturally be of the same order, electroweak symmetry breaking can be obtained quite naturally in this model.

V. CONCLUSIONS

In the MSSM the upper bound on the mass of the lightest neutral Higgs boson is about 130 GeV. We have constructed simple, realistic supersymmetric models which allow the mass of this particle to significantly exceed 130 GeV while maintaining the MSSM prediction for gauge coupling unification. These models are based on warped 5D versions of the NMSSM (or related theories). The point is that by localizing the NMSSM singlet on the IR brane it is possible for the couplings of this field to become large at the IR cutoff without affecting the prediction of gauge coupling unification. These can give a large contribution to the quartic coupling of the Higgs field, thereby raising the bound on the mass of the lightest neutral Higgs boson.

Several versions of these models are possible, differing primarily in the location of the quark and lepton fields and in the pattern of breaking of the unified gauge symmetry. The exact bound on the Higgs boson mass is somewhat different in each of these models. In this paper we have concentrated primarily on two particularly simple cases. In the first, the quark and lepton fields are localized on the UV brane, where the unified symmetry is broken. There are (at least) two sets of Higgs doublets—one localized on the IR brane receiving a large quartic coupling from the NMSSM potential on the IR brane, the other propagating the bulk having the Yukawa couplings to the quarks and leptons. The Higgs doublets responsible for electroweak symmetry breaking are linear combinations of these two sets, thus having both the Yukawa couplings and a large quartic coupling. The running of the top Yukawa coupling, however, is nonstandard above the IR scale and it limits the maximum value of the bound we can obtain in this model. In the second model, the quark and lepton fields are localized on the IR brane, while the unified symmetry is broken both on the UV and IR branes. In this case there is no issue of the nonstandard running of the top Yukawa coupling so that we can obtain the maximal value of the bound on the Higgs mass, which can be as large as 300 GeV.

Through the AdS/CFT correspondence, our theory is related to purely 4D theories in which some (or all) of the Higgs fields are composites of some strong interaction, which is nearly conformal above the scale of the IR cutoff. In this sense our theory shares certain features with the model considered in [28]. In our model, however, the prediction of gauge coupling unification is “automatically” the same as the MSSM, which is not the case in the model of [28]. The issues of raising the Higgs mass bound and the prediction of the gauge couplings are also considered in [29]. In that model, however, in contrast to

ours, the Higgs doublets are mainly elementary while the singlet is composite.

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