# Higgs descendants

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# Higgs Descendants

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We define a Higgs descendant $\chi$ to be a particle beyond the standard model whose mass arises predominantly from the vacuum expectation value of the Higgs boson. Higgs descendants arise naturally from new physics whose intrinsic mass scale is unrelated to the electroweak scale. The coupling of $\chi$ to the Higgs boson is fixed by the mass and spin of $\chi$, yielding a highly predictive setup in which there may be substantial modifications to the properties of the Higgs boson. For example, if the decay of the Higgs boson to $\chi$ is kinematically allowed, then this branching ratio is largely determined. Depending on the stability of $\chi$, Higgs decays may result in a variety of possible visible or invisible final states. Alternatively, loops of $\chi$ may affect Higgs boson production or its decays to standard model particles. If $\chi$ is stable dark matter, then the mandatory coupling between $\chi$ and the Higgs boson gives a lower bound on the direct detection cross section as a function of the $\chi$ mass. We also present a number of explicit models which are examples of Higgs descendants. Finally, we comment on Higgs descendants in the context of the excesses near 125 GeV recently observed at ATLAS and CMS.

## I. INTRODUCTION

A primary goal of the LHC is to reveal the fundamental dynamics underlying electroweak symmetry breaking. The characteristic mass scale of the standard model is the vacuum expectation value (VEV) of the Higgs field, $v = 246$ GeV, which is radiatively unstable. This fact has been the central driving force for exploring new theories beyond the standard model at the TeV scale.

Theories which address the electroweak hierarchy problem universally introduce dimensionful parameters from which $v$ originates. For example, in the case of supersymmetry [1] the $\mu$ term and soft masses determine the value of $v$. Likewise, in theories with a pseudo Nambu-Goldstone Higgs boson [2], the scale $v$ descends a loop down from the symmetry breaking decay constant $f$. In such instances the proximity of new mass scales to $v$ is *required*, and thus not a coincidence.

Conversely, for any new physics unrelated to the origin of electroweak symmetry breaking, associated dimensionful parameters have, a priori, no reason to be near $v$. If these mass scales lie far above the electroweak scale, then the associated particles are kinematically inaccessible to experiments like the LHC, and we can integrate them out. Hence, in this limit our only hope of probing such new physics is from higher dimension operators, e.g. as in the case of the Weinberg operator for small neutrino masses [3].

If, on the other hand, the new mass scales are very tiny then the associated states will be, naively, quite light. This naive intuition, however, fails if the new physics is coupled to the standard model. In this case there will in general be induced interactions between these new states and the Higgs field. After electroweak symmetry breaking, such couplings provide the main contribution to the masses of these new states, which are necessarily proportional to some power of $v$. We will refer to this broad class of new states as Higgs descendants, and as we will see they can have a drastic impact on Higgs physics even though they do little to address the hierarchy problem.

More concretely, let us define a Higgs descendant as an additional particle beyond the standard model, $\chi$, whose mass obeys the property that

$$v \to 0 \quad \Rightarrow \quad m_\chi \to 0. \tag{1}$$

Note that this condition does not hold typically for new states associated with the sectors addressing the hierarchy problem. For instance, in supersymmetry the superpartner masses arise from soft supersymmetry breaking parameters, so they do not necessarily vanish as $v \to 0$. Likewise, any particle with a bare mass is not a Higgs descendant.

Still, Higgs descendants represent large classes of theories that may appear in varying contexts. Several examples of Higgs descendants can already be found in the literature. For instance, a number of authors have considered minimal singlet dark matter coupled via the so-called Higgs portal [4, 5]. In many cases, the singlet is given an intrinsic mass scale which may be fixed by other considerations, e.g. a thermal relic abundance [6]. Restricting such theories to forbid/suppress explicit mass terms for the dark matter yields a Higgs descendant [7]. Likewise, fourth generation Dirac neutrinos [8] and certain hidden sector models [9, 10] are also examples of Higgs descendants.

The condition of Eq. (1) implies that $m_\chi$ can be expanded in powers of $v$, so for example

$$m_\chi = \lambda v^n + \dots, \tag{2}$$

for a fermionic Higgs descendant $\chi$. Here, $n > 0$, and $\lambda$ parametrizes our ignorance of physics coupling the Higgs field to the Higgs descendant. For weakly coupled theories, $n$ is integer, but in general it need not be [11]. The ellipses denote higher order terms in $v$, which we will ignore throughout. Higgs descendant theories possess a highly restricted phenomenology because the interactions relevant to many physical processes are essentially fixed

by $m_\chi$ alone. In particular, by replacing the Higgs VEV as

$$v \to v + h, \qquad (3)$$

where $h$ is the propagating Higgs boson field, Eq. (2) leads to a Lagrangian of the form

$$\mathcal{L} = -m_\chi \left(1 + \frac{nh}{v}\right) \bar\chi\chi + \dots, \qquad (4)$$

where the ellipses denote terms higher order in the Higgs field and we have considered a Dirac fermion $\chi$ for concreteness.

For a scalar Higgs descendant $\chi$, we define the index $n$ according to

$$m_\chi^2 = \lambda v^n + \dots, \qquad (5)$$

reflecting the dimension of the $\chi$ mass term in the Lagrangian. Equation (4) then applies equally for a complex scalar $\chi$ only by making $m_\chi \to m_\chi^2$. The generalization is obvious for other spins and real representations. Note that throughout, $n$ is defined by Eq. (2) and Eq. (5) for fermionic and bosonic $\chi$, respectively.

The organization of our paper is as follows. In Section II, we discuss how a Higgs descendant can have significant implications for Higgs phenomenology at the LHC. These arise either as new decay modes or as modifications of the standard model production and decay modes of the Higgs boson. We discuss the case in which the Higgs descendants are stable or unstable decaying into standard model particles, leading to distinct signatures for Higgs boson decays. In Section III, we consider the case in which a Higgs descendant is stable and comprises the dark matter of the universe. In this case, Eq. (4) provides a lower bound, modulo unnatural cancellations, on the dark matter-nucleon scattering cross section relevant for direct detection experiments, as a function of $m_\chi$. Correlations between dark matter and LHC physics would then provide a powerful probe of the underlying theory. In Section IV, we present simple, explicit models of Higgs descendants. Finally, we conclude in Section V.

## II. HIGGS BOSON PROPERTIES

The existence of a Higgs descendant $\chi$ leads to the mandatory coupling in Eq. (4). Here we discuss the effect of this coupling on Higgs physics. Throughout this paper, we will limit ourselves to the case of a single Higgs doublet.

### A. Light Higgs Descendants

Consider a scenario of light Higgs descendants, defined as $m_\chi < m_h/2$. In this case the Higgs boson has a new decay channel, $h \to \chi\chi$, with a branching ratio fixed by $m_h$, $m_\chi$, and $n$, which is independent of the gauge quantum numbers of $\chi$ (modulo multiplicity factors).[1]

The partial decay rate of the Higgs boson to standard model particles can be found in Ref. [12], while the partial decay rate to $\chi$ is given by

$$\Gamma_f = \frac{m_h}{8\pi}\left(\frac{nm_\chi}{v}\right)^2 \left(1 - \frac{4m_\chi^2}{m_h^2}\right)^{3/2}, \qquad (6)$$

for a Dirac fermion $\chi$ and by

$$\Gamma_s = \frac{m_h}{16\pi}\left(\frac{m_\chi}{m_h}\right)^2 \left(\frac{nm_\chi}{v}\right)^2 \left(1 - \frac{4m_\chi^2}{m_h^2}\right)^{1/2}, \qquad (7)$$

for a complex scalar $\chi$. For a Majorana fermion or a real scalar, one makes the replacements $\Gamma_f \to \Gamma_f/2$ and $\Gamma_s \to \Gamma_s/2$, respectively. The additional factor of $m_\chi^2/m_h^2$ in Eq. (7) relative to Eq. (6) will produce quantitatively different physics for fermionic versus scalar $\chi$, especially for small $m_\chi$.

The branching ratio $\mathrm{BR}(h \to \chi\chi)$ is shown in Fig. 1 for a Dirac fermion $\chi$ as a function of $m_\chi$ and $m_h$ for $n = 1$ and 2. One expects that for $m_h \lesssim 2m_W$, the Higgs branching ratio into a Higgs descendant becomes important when $nm_\chi$ is of order $m_b$, the bottom quark mass, and this is indeed the case. For example, $\mathrm{BR}(h \to \chi\chi) \gtrsim 50\%$ for a light Higgs boson when $m_\chi \gtrsim 6$ GeV ($m_\chi \gtrsim 3$ GeV) for $n = 1$ ($n = 2$). For a heavy Higgs boson above the $WW$ threshold, decays into the Higgs descendant can be non-negligible only in small regions of parameter space.

If $\chi$ is a scalar, then the story changes quantitatively. In this case, $\chi$ interacts with the Higgs boson with a dimensionful coupling $nm_\chi^2/v$, rather than $nm_\chi/v$. This induces an additional factor of $m_\chi^2/m_h^2$ which suppresses the Higgs boson decay rate to Higgs descendants, especially for small $m_\chi$. This trend is verified in Fig. 2, which depicts $\mathrm{BR}(h \to \chi\chi)$ for complex scalar $\chi$. We find that $\mathrm{BR}(h \to \chi\chi) \gtrsim 50\%$ for a light Higgs boson only when $m_\chi \gtrsim 30$ GeV ($m_\chi \gtrsim 20$ GeV) for $n = 1$ ($n = 2$). Decays into scalar Higgs descendants never dominate for a heavy Higgs boson above the $WW$ threshold.

What occurs after the Higgs decays to a pair of $\chi$ particles is somewhat more model dependent. Hence, the resulting experimental signatures are as well. The interactions in Eq. (4) preserve a $\mathbb{Z}_2$ symmetry which, if unbroken, ensures the stability of $\chi$. Consider first the case where this $\mathbb{Z}_2$ is sufficiently exact that $\chi$ is stable on collider time scales. In this case, the signatures depend largely on the standard model gauge quantum numbers of $\chi$, which we now consider cases by case.

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[1] If $n$ is odd, it might naively be thought that $\chi$ must carry standard model charges, but this is not true. The mass of $\chi$ may arise from a singlet field $s$ whose VEV is induced by the Higgs VEV, i.e. $\langle s \rangle \propto v$.
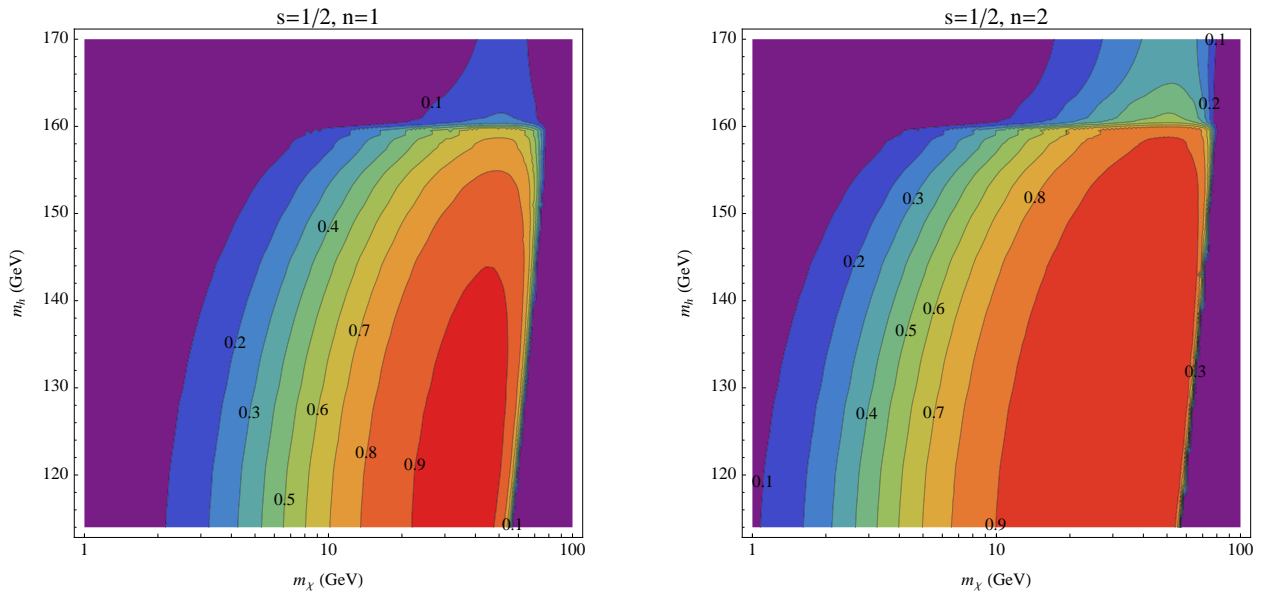
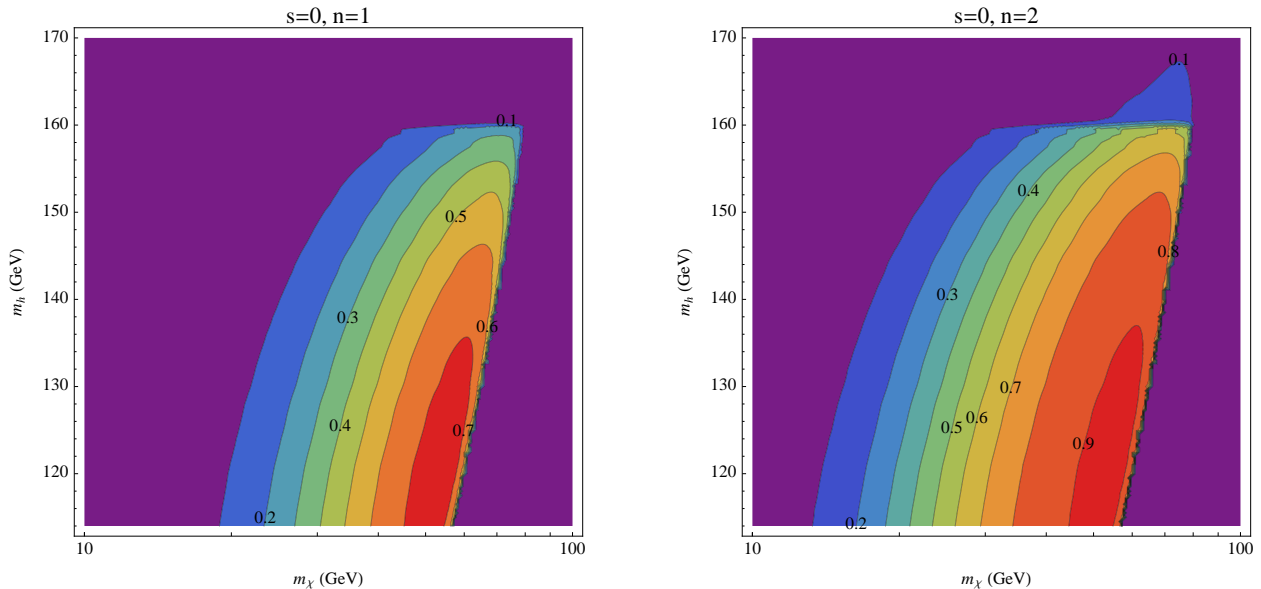FIG. 1: Branching ratio of $h \to \chi\chi$ for a (Dirac) fermionic $\chi$ with $n = 1$ (left) and $n = 2$ (right).



FIG. 2: Branching ratio of $h \to \chi\chi$ for a (complex) scalar $\chi$ with $n = 1$ (left) and $n = 2$ (right).

### 1. Stable Case

If $\chi$ is colored it will hadronize, yielding highly ioniz-ing, sometimes intermittent tracks caused by long-lived, charged hadronic states.[2] This possibility, however, is

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[2] For discussions of such signatures in other contexts, see e.g. Ref. [13]. Note also that the Higgs decay widths in this case are $N$ times those in Eqs. (6) and (7), where $N = 3$ and 8 for a color triplet and octet $\chi$, respectively.

strongly constrained by direct production of $\chi$ at the LHC [14], and is not viable unless the hadronization leads mostly to neutral bound states.

Next, let us consider the scenario in which the Higgs descendant is uncolored, but carries electroweak charge. For example, if the Higgs descendant has the quantum numbers of a lepton, then $\chi = (\chi^0, \chi^+)$. In this case the Higgs boson may decay into the charged components $\chi^+$, possibly as well as the neutral components $\chi^0$, depend-ing on the quantum numbers of the $\chi/\bar{\chi}$ multiplets. The Higgs descendant then may lead to signatures discussed in Ref. [15] if the mass splitting $\Delta m \equiv m_{\chi^+} - m_{\chi^0}$ is suf-

ficiently small—for $\Delta m \lesssim m_\pi$ it leads to stable charged tracks, while for $m_\pi \lesssim \Delta m \lesssim 200$ MeV to short tracks of $\chi^+$ together with soft pions from $\chi^+ \to \chi^0 \pi^+$. If $\Delta m$ is larger (but not much larger than $\simeq$ GeV), then both $h \to \chi^+ \chi^-$ and $h \to \chi^0 \chi^0$ are recognized only as invisible Higgs boson decays at the LHC (unless boost factors for $\chi^+$ are large due to $m_\chi \ll m_h$), with the widths larger than those in Eqs. (6) and (7) due to an appropriate multiplicity factor. For $\Delta m >$ GeV, decay products of $\chi^+$ may be directly tagged.

Finally, if $\chi$ is not charged under any standard model gauge interactions, then $h \to \chi\chi$ contributes to the invisible Higgs decay width. (For recent discussions on invisibly decaying Higgs bosons, see e.g. Ref. [16].) If exactly stable and neutral, $\chi$ is a possible dark matter candidate, which we discuss in depth in Sec. III.

### 2. Unstable Case

In general, the $\mathbb{Z}_2$ symmetry preserved by Eq. (4) may be broken if there are additional couplings of $\chi$ to standard model particles. These interactions permit the Higgs descendant to decay visibly inside the collider, as discussed in a similarly general context in Ref. [17]. For example, for a singlet fermionic $\chi$, consider the interaction $\chi\mathcal{O}$, where

$$\mathcal{O} = \{u^c d^c d^c, \ell\ell e^c, q\ell d^c, q\ell^\dagger u^c, u^c d^{c\dagger} e^c, \ell h\}. \quad (8)$$

Here, $\mathcal{O}$ is constructed from all leading singlet fermionic standard model composites. Each of these operators affects the phenomenology in a distinct way.

For $\mathcal{O} = u^c d^c d^c$ each Higgs descendant decays to three jets, yielding a striking six jet final state from Higgs boson decays. In this scenario, decays of the Higgs boson may be buried in soft jets and thus difficult to discern at the LHC. This operator also induces low energy mixing between $\chi$ and standard model baryons. For $\mathcal{O} = \ell\ell e^c$, the Higgs descendant decays via di-lepton plus missing energy. For $\mathcal{O} = q\ell d^c, q\ell^\dagger u^c, u^c d^{c\dagger} e^c$, each Higgs descendant decays to di-jet plus lepton (or di-jet plus missing energy for the first two operators). For $\mathcal{O} = \ell h$, the Higgs descendant decays to di-jet or di-lepton plus missing energy, or di-jet plus lepton. However, this operator is not ideal because it generates a large Dirac mass between $\chi$ and neutrinos, which is difficult to reconcile with neutrino oscillation experiments and cosmological constraints.

Lastly, consider the case of a singlet scalar $\chi$, which can couple via $\chi\mathcal{O}$ where

$$\mathcal{O} = \{qhu^c, qh^\dagger d^c, \ell h^\dagger e^c, F_{\mu\nu}F^{\mu\nu}, G_{\mu\nu}G^{\mu\nu}\}. \quad (9)$$

Depending on the operator, $\chi$ will decay to di-jet, di-lepton, or di-gamma.

It is straightforward to apply similar arguments to $\chi$ which is not a standard model singlet, but we will not do so here.

### B. Heavy Higgs Descendants

Let us now consider the case of heavy Higgs descendants, where $m_\chi > m_h/2$. In this case, the Higgs boson is kinematically forbidden from decaying to Higgs descendants. One might think naively that the properties of the Higgs boson are very similar to that of the standard model. However, in certain cases $\chi$ is charged under the standard model gauge group, in which case it will typically influence Higgs boson production via $gg \to h$ as well as the decay $h \to \gamma\gamma$. Moreover, as we will see, the contributions to these processes from Higgs descendants are fixed by the choice of $n$ and the spin and charges of $\chi$.

The coupling of the Higgs boson to the gluon is

$$\begin{aligned}
\mathcal{L}_{hgg} &= \frac{\alpha_s}{12\pi} \frac{h}{v} \left( \sum_i c_i t_i N_i \frac{\partial \log m_i(v)}{\log v} \right) G^a_{\mu\nu} G^{a\mu\nu} \\
&= \frac{\alpha_s}{12\pi} \frac{h}{v} \left( 1 + c_\chi t_\chi N_\chi n \right) G^a_{\mu\nu} G^{a\mu\nu}, \quad (10)
\end{aligned}$$

where $i$ labels heavy species which provide threshold corrections to the beta function of QCD. Here $c_i = 2$ for Dirac fermions and $c_i = 1/2$ for complex scalars, $t_i$ is the Dynkin index of the multiplet, and $N_i$ is the multiplicity. In the second line we have plugged in for $i = \chi$ and the top quark, where $\chi$ is colored. We see that the Higgs boson coupling to the gluon is fixed by a set of discrete quantum numbers of the $\chi$ field. For example, the fourth generation quarks have $\{c_\chi, t_\chi, N_\chi, n\} = \{2, 1/2, 2, 1\}$, enhancing the amplitude for Higgs production through gluon fusion by a factor of three and thus the production cross section by a factor of nine.

A similar formula for the coupling of Higgs boson to the photon can be derived from the beta function of QED:

$$\begin{aligned}
\mathcal{L}_{h\gamma\gamma} &= \frac{\alpha}{12\pi} \frac{h}{v} \left( -\frac{21}{2} + \sum_i c_i q_i^2 N_i \frac{\partial \log m_i(v)}{\log v} \right) F_{\mu\nu} F^{\mu\nu} \\
&= \frac{\alpha}{12\pi} \frac{h}{v} \left( -\frac{21}{2} + \frac{8}{3} + c_\chi q_\chi^2 N_\chi n \right) F_{\mu\nu} F^{\mu\nu}, \quad (11)
\end{aligned}$$

where the Dynkin index $t_i$ is replaced with the electric charge $q_i^2$, and the term of $-21/2$ is the contribution to this coupling from a $W$ boson loop. As before, we have plugged in for $i = \chi$ and the top quark in the second line, assuming an electrically charged Higgs descendant. As is well-known, the top quark and $W$ loop contributions to the Higgs coupling to photons destructively interfere [12]. Similarly, loops of Higgs descendants also cancel against the $W$ loop contribution, since the mass of a Higgs descendant always grows with the Higgs VEV. For example, this occurs for a fourth generation lepton, which contributes $\{c_\chi, q_\chi, N_\chi, n\} = \{2, 1, 1, 1\}$.

The defining relation of our scenario, Eq. (2), implies that the mass of a Higgs descendant cannot be arbitrarily large, $m_\chi \lesssim 1$ TeV, so it is subject to direct search limits from the LHC. We now consider those limits. While there

are no dedicated LHC searches for states with arbitrary charges, we can get rough estimates for the limits from similar searches.

If $\chi$ is long-lived on collider time scales, then LHC searches for stable colored or electrically charged particles apply. These mass limits are quite stringent for gluino-like ($\gtrsim 1$ TeV) and squark-like ($\gtrsim 700$ GeV) states, but relatively weak for slepton-like ($\gtrsim 200$ GeV) states [14]. We can thus conclude that long-lived Higgs descendants are substantially constrained if colored, but are quite viable otherwise.

On the other hand, if $\chi$ decays promptly then limits depend sensitively on its charges and decay modes. For instance, consider $\chi$ which is colored and promptly decays. For $\chi$ decaying to standard model quarks plus one or more leptons, the decay topology and thus the LHC limit are similar to that of top and bottom partner quarks ($\gtrsim 600$ GeV) [18]. Meanwhile, for $\chi$ decaying to standard model quarks plus missing energy, supersymmetry searches apply. Using the simplified model analysis of Ref. [19], one can estimate the LHC limit for a single squark-like state ($\gtrsim 300$ GeV), although this bound can be eliminated altogether if the mass of the missing energy particle is sufficiently large. If, on the other hand, $\chi$ decays entirely to jets with no missing energy, as is the case in many $R$-parity violating models, bounds will be substantially weaker.

Next, let us consider the case of color-neutral, electroweak charged $\chi$ which promptly decays. In this scenario, limits are far weaker due to the smaller production cross section. For chargino-like states, one can apply constraints from LHC supersymmetry searches ($\gtrsim 300$ GeV) [20] although these bounds make very specific assumptions about the cascade decay—in particular these searches are driven by the presence of intermediate sleptons which provide additional leptons. In general, the limits on the minimal chargino-neutralino system are very weak at LHC, and the dominant bounds still come from LEP ($\gtrsim 100$ GeV).

Let us now determine the effect of Eq. (10) and Eq. (11) on Higgs boson phenomenology. To do so we define a function $r[x]$ to be the ratio of a quantity $x$ in a particular theory divided by the value of $x$ in the standard model. In Fig. 3 we have plotted $r[\mathrm{BR}(h \to \gamma\gamma)]$ for a color singlet $\chi$ with electric charge equal to $q_\chi$. As expected, the branching ratio to di-gamma diminishes for small values of $q_\chi$ due to the destructive interference with the $W$ boson. For sufficiently large $q_\chi$, however, the contribution from $\chi$ dominates and the di-gamma branching ratio begins to increase. As discussed, a color singlet $\chi$ is relatively unconstrained by direct search limits and can be coupled perturbatively to the Higgs boson. In Fig. 4, we have plotted $r[\sigma(gg \to h)\mathrm{BR}(h \to \gamma\gamma)]$ for a color triplet $\chi$ with electric charge $q_\chi$. Here the rate of di-gamma events from Higgs boson decays is substantially increased at low $q_\chi$, but can again cancel at larger values. We emphasize, however, that because of direct search limits on colored particles, much of this param-
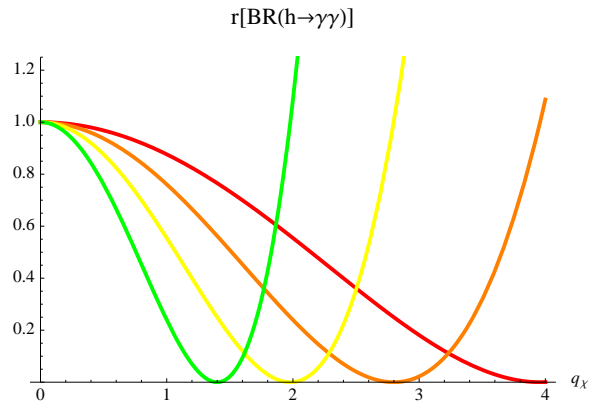


FIG. 3: Ratio of $\mathrm{BR}(h \to \gamma\gamma)$ versus its standard model value for a color singlet $\chi$ of electric charge $q_\chi$. Here, {red, orange} correspond to a complex scalar with $n = \{1, 2\}$ and {yellow, green} correspond to a Dirac fermion with $n = \{1, 2\}$.
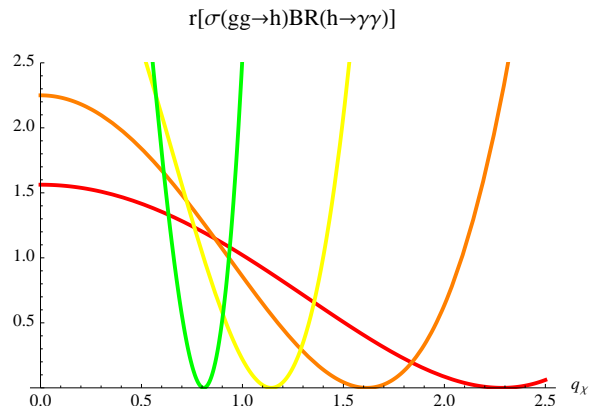


FIG. 4: Ratio of $\sigma(gg \to h)\mathrm{BR}(h \to \gamma\gamma)$ versus its standard model value for a color triplet $\chi$ of electric charge $q_\chi$. The colors are as in Fig. 3.

eter space is disfavored, without hiding $\chi$ by making it decay into multiple jets or invoking non-perturbative dynamics to sufficiently lift the mass of $\chi$.

## III. HIGGS DESCENDANT DARK MATTER

If a Higgs descendant $\chi$ contains a neutral and stable component, then this can comprise the dark matter of the universe. In this section, we consider this possibility, especially its implications for direct detection experiments. In what follows, we do not impose any constraints from demanding a thermal relic abundance of $\chi$ particles through freeze-out through the interaction in Eq. (4). As noted in Ref. [21], for example, there is much leeway in generating dark matter of this type through non-thermal methods.

As a consequence of Eq. (4), $\chi$ has a mandatory cou-

pling to the Higgs boson. Since the Higgs boson couples to quarks and gluons, one in turn should expect an irreducible spin-independent scattering cross section of $\chi$ against a target nucleus via $t$-channel Higgs boson exchange. For a Dirac fermion $\chi$, the cross section is

$$\sigma = \frac{\mu^2}{\pi}\left(\frac{nm_\chi}{v}\right)^2\frac{1}{m_h^4}\left(Zg_{hpp}+(A-Z)g_{hnn}\right)^2, \quad (12)$$

where $\mu$ is the dark matter-nucleus reduced mass, $m_h$ is the Higgs boson mass, $Z$ and $A$ are the atomic number and weight of the target nucleus, and $g_{hNN}$ for $N=p,n$ are the coupling of the Higgs boson to the proton and neutron:

$$g_{hNN} = \frac{m_N}{v}\sum_q f_q^N. \quad (13)$$

The scattering cross section for complex scalar $\chi$ is one quarter that of a Dirac fermion.

To derive the dark matter-nucleon spin-independent cross section $\sigma_{\rm SI}$, we set $A = Z = 1$ and $\mu = 1$ GeV. Using $\sum_q f_q^N \simeq 0.30 \pm 0.015$, compiled in Ref. [6] based on Ref. [22], we obtain

$$\sigma_{\rm SI} \simeq 2.5\times10^{-45}\ {\rm cm}^2\left(\frac{nm_\chi}{10\ {\rm GeV}}\right)^2\left(\frac{114\ {\rm GeV}}{m_h}\right)^4, \quad (14)$$

for Dirac fermion dark matter. This gives a lower bound on the dark matter detection cross section; unless there is a substantial cancellation between the Higgs-exchange and other processes, we expect that the true dark matter-nucleon cross section is comparable to or larger than the value given in Eq. (14).

The above result can now be compared with predictions for the invisible Higgs boson decay rate: $h \to \chi\chi$, where $\chi$ is neutral and stable component of the Higgs descendant. In Figs. 5 and 6, we present contour plots of $BR(h \to \chi\chi)$ in the $(m_\chi, \sigma_{\rm SI})$ planes for complex scalar and Dirac fermion dark matter, respectively. Assuming that the dark matter-nucleon cross section saturates the lower bound in Eq. (14), each point in the $(m_\chi, \sigma_{\rm SI})$ plane corresponds to a fixed value of $m_h$ for a given $n$, and thus to a definite value of $BR(h \to \chi\chi)$. Using this, we have shown shaded bands corresponding to the Higgs boson mass range 114 GeV $< m_h <$ 145 GeV for $n = 1, 2$ in each figure. Within these bands we have also drawn labeled contours, showing $BR(h \to \chi\chi)$. Finally, the dashed black line in each figure indicates the most recent exclusion curve from the XENON100 dark matter detection experiment [23]. From Fig. 5, we see that if an invisible branching ratio of the Higgs boson greater than $\sim 70\%$ is measured, then complex scalar dark matter of this type will be excluded (up to astrophysical uncertainties associated with the dark matter detection constraint). On the other hand, from Fig. 6 it is clear that for Dirac fermion dark matter, a mostly invisibly decaying Higgs boson can easily be accommodated.
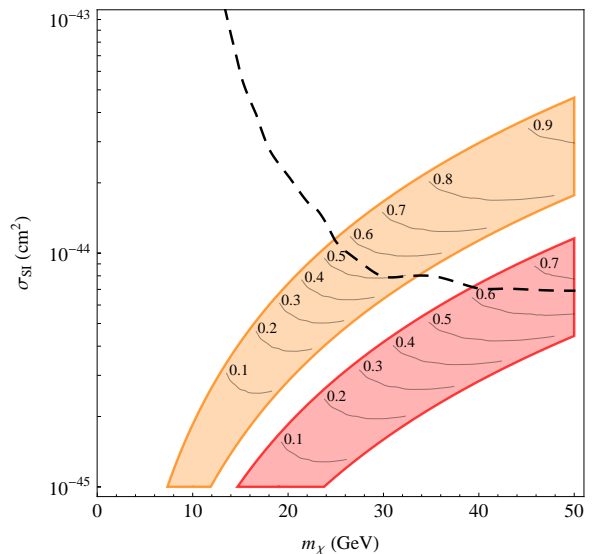


FIG. 5: Plot of $\sigma_{\rm SI}$ versus $m_\chi$ for complex scalar dark matter. The dashed black line depicts the constraint from XENON100. The {red, yellow} bands indicate the range corresponding to 114 GeV $< m_h <$ 145 GeV for $n = \{1, 2\}$, respectively, while the labeled contours indicate the value of $BR(h \to \chi\chi)$.
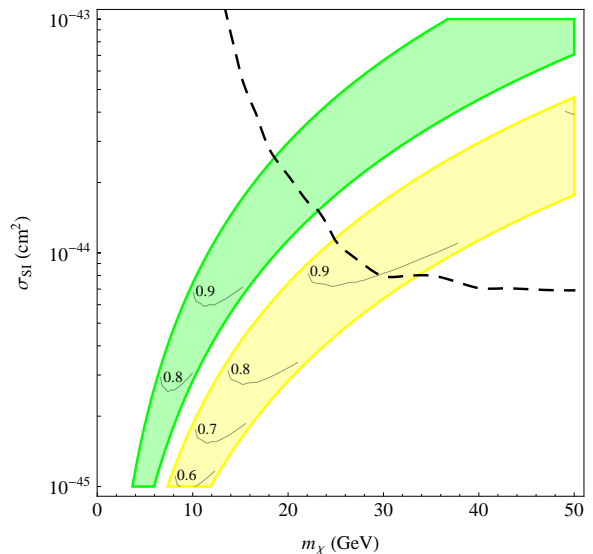


FIG. 6: Same as Fig. 5 except for Dirac fermion dark matter. The {yellow, green} bands indicate the range corresponding to 114 GeV $< m_h <$ 145 GeV for $n = \{1, 2\}$, respectively.

## IV. EXPLICIT MODELS

A number of prominent examples of Higgs descendants have been studied in the existing literature, albeit in varying contexts. Perhaps the most obvious among these theories is singlet dark matter coupled via the so-called

Higgs portal [4, 6, 7]:

$$\mathcal{L} = -\frac{\lambda}{4}|h|^2|S|^2. \qquad (15)$$

Because this model is rather phenomenological in nature, naturalness considerations are disregarded and a bare mass term for the gauge singlet, $m^2|S|^2$, is simply not included.[3] The analyses of Secs. II and III apply more or less verbatim to this model, whose properties are correctly characterized by Figs. 2 and 5.

A less obvious example of a Higgs descendant is a fourth generation lepton, consisting of $L, E^c$ and $N^c$ with the interaction

$$\mathcal{L} = \lambda L h N^c + \kappa L h^\dagger E^c, \qquad (16)$$

where $L = (N, E)$ is a left-handed doublet.[4] The right-handed neutrino, $N^c$, is a gauge singlet, so in general it can have a Majorana mass term

$$\mathcal{L} = -\frac{m}{2}N^c N^c. \qquad (17)$$

However, as discussed earlier, the value of $m$ has no reason to be near the electroweak scale, so we expect it not to be very close to $v$.

If $m$ is sufficiently larger than $v$, then $N^c$ should be integrated out, yielding a higher dimension Majorana mass for $N$ given by $m_N = \lambda^2 v^2/m$. Since $N$ acquires a mass solely from electroweak symmetry breaking, it is an example of a Higgs descendant, as expected from our general considerations. Unfortunately, in the limit $m \gg v$ one expects $m_N < m_Z/2$, so this theory is disfavored from $Z$ pole measurements.

Alternatively, it may be that $m \ll v$, in which case $N^c$ remains in the spectrum and acquires a Dirac mass with $N$ such that $m_N = \lambda v$. In this case $(N, N^c)$ is a Higgs descendant and Figs. 1 and 6 pertain. For $m_N > m_Z/2$ this theory trivially evades $Z$ pole constraints, but having $(N, N^c)$ dark matter may be difficult. This is because Fig. 6 implies $m_h \gtrsim 175$ GeV, in which case there is no way to avoid the bound from standard model Higgs search [24], given that $\mathrm{BR}(h \to WW^*)$ cannot be suppressed; see Fig. 1. The $(N, N^c)$ dark matter is possible if the XENON100 constraint is a factor of 2 weaker than that depicted in Fig. 6, due e.g. to astrophysical uncertainties. In this case, the direct detection constraint requires only $m_h \gtrsim 150$ GeV, so that significant depletion of $\mathrm{BR}(h \to WW^*)$ is possible for $m_N < m_h/2$.

Let us now consider our final example. Suppose there is a supersymmetric "hidden sector" which has

a $U(1)_X$ gauge field kinetically mixing with the hypercharge $U(1)_Y$ of the standard model [25]:

$$\mathcal{L} = \frac{\epsilon}{2}\int d^2\theta \, \mathcal{W}_Y \mathcal{W}_X \supset \epsilon D_Y D_X. \qquad (18)$$

The mixed $D$-term then produces a scalar potential of the form

$$V = \frac{1}{2}\left\{g_X(|X|^2 - |X^c|^2) + \xi\right\}^2, \qquad (19)$$

$$\xi = \epsilon D_Y = -\frac{\epsilon g_Y}{2}|h|^2. \qquad (20)$$

Here, $X$ and $X^c$ are hidden sector chiral superfields charged under $U(1)_X$. For simplicity, we have taken the decoupling limit where $h$ is really the up-type Higgs boson at $\tan\beta = \infty$. After electroweak symmetry breaking, $\xi$ acquires a VEV, triggering spontaneous symmetry breaking in the hidden sector proportional to the order parameter

$$\langle X \rangle = \sqrt{\frac{|\xi|}{g_X}}, \qquad (21)$$

where we have assumed $\epsilon > 0$. Since $\langle X \rangle$ breaks the $U(1)_X$ gauge symmetry, the real and imaginary components of $X$ are eaten to become the longitudinal and radial modes of the massive vector supermultiplet, $V_X$. The mass of this supermultiplet is given by $m_{V_X} = \sqrt{2g_X|\xi|} \propto v$, so that $V_X$ is a Higgs descendant and our arguments apply to all its components. (These components split in mass after supersymmetry breaking.)

## V.  DISCUSSION

In this paper we have analyzed a broad class of theories in which a new particle beyond the standard model, $\chi$, acquires its mass predominantly from the VEV of the Higgs field, $v$, rather than any intrinsic mass scale. Such a particle, which we called a Higgs descendant, can arise naturally from any new sector whose intrinsic mass scales does not coincide with the electroweak scale.[5] Because the couplings of Higgs descendants are highly constrained, as given in Eq. (4), the physics is dictated essentially by the mass and spin of $\chi$. As we have seen, both for light and heavy Higgs descendants there can be substantial modifications of the production and decay of the Higgs boson. In the case where $\chi$ is also a dark matter particle, this class of theories predict a minimum spin-independent direct detection cross section that could be probed in experiments in the near future.

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[3] Such a model can be made natural by the inclusion of supersymmetry.

[4] The standard model gauge anomalies can be canceled by introducing, e.g., a mirror lepton generation $(L^c, E, N)$ or a fourth generation quark $(Q, U^c, D^c)$.

[5] The QCD scale is another scale in the standard model, which is a priori independent of the electroweak symmetry breaking scale. In principle, one may consider "QCD descendant" by coupling new physics to a QCD condensate.
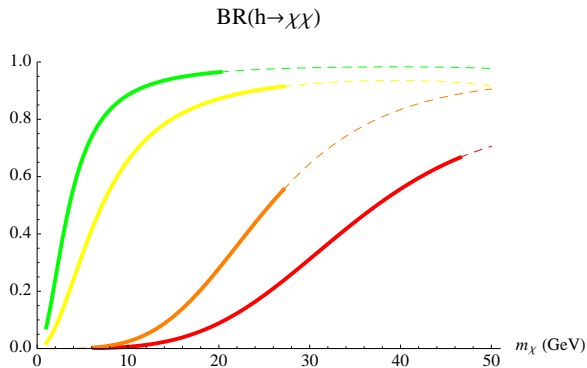
FIG. 7: Plot of BR($h \to \gamma\gamma$) versus dark matter mass $m_\chi$ for $m_h = 125$ GeV. Here, {red, orange} corresponds to a complex scalar with $n = \{1, 2\}$ and {yellow, green} to a Dirac fermion with $n = \{1, 2\}$. Dashed regions are excluded by XENON100.

While our discussion has been limited to the case of a single Higgs doublet theory, our analysis can be straightforwardly extended to a two Higgs doublet model simply by expanding $m_\chi$ in a power series in $v_u$ and $v_d$. However, this class of Higgs descendants becomes less predictive, due to a proliferation of free parameters coming from the multi-variate expansion of $m_\chi$ as well as Higgs boson mass and VEV mixing angles such as $\tan\alpha$ and $\tan\beta$. Our analysis, however, applies without modification in the decoupling regime.

Likewise, our discussion can also be extended to include supersymmetry. In many such theories, the inter-

action of Eq. (4) is accompanied by a supersymmetric analog in which the Higgs boson is replaced by the Higgsino. Hence, any given value of $m_\chi$ suggests a minimal coupling of the Higgsino to $\chi$ and its superpartner. We leave an investigation of these possible Higgsino descendants for future work.

Lastly, let us briefly comment on the recently observed excesses seen at ATLAS and CMS in the di-gamma and $ZZ$ channels [26]. Interpreted as a signal from the decay of the Higgs boson, this suggests a mass $m_h \simeq 125$ GeV, although these experiments obviously cannot yet make definite statements about the detailed properties of such a Higgs particle. As such, it is essential that further experimental analyses is applied to determine more precisely these properties. Fixing $m_h = 125$ GeV, any Higgs descendant theory has a parameter space which is even more restricted than that discussed in previous sections. In Fig. 7 we have plotted the value of BR($h \to \chi\chi$) for $n = 1, 2$ and for Dirac fermion and complex scalar $\chi$. The solid (dashed) portions of the curve correspond to regions in parameter space which are (dis)allowed by XENON100, if $\chi$ composes all of the dark matter.

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