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Hidden spin-3/2 field in the standard model

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Abstract Here we show that a massive spin-3/2 field can hide in the SM spectrum in a way revealing itself only virtually. We study collider signatures and loop effects of this field, and determine its role in Higgs inflation and its potential as dark matter. We show that this spin-3/2 field has a rich linear collider phenomenology and motivates consideration of a neutrino–Higgs collider. We also show that the study of Higgs inflation, dark matter and dark energy can reveal more about the neutrino and dark sector.

1 Introduction

The standard model (SM) of strong and electroweak interactions, spectrally completed by the discovery of its Higgs boson at the LHC [1], seems to be the model of physics at Fermi energies. It does so because various experiments have revealed so far no new particles beyond the SM spectrum. There is, however, at least the dark matter (DM), which requires new particles beyond the SM. Physically, therefore, we must use every opportunity to understand where those new particles can hide, if there are any.

In the present work we study a massive spin-3/2 field hidden in the SM spectrum. This higher-spin field, described by the Rarita–Schwinger equations [2–4], has to obey certain constraints to have correct degrees of freedom when it is on the physical shell. At the renormalizable level, it can couple to the SM matter via only the neutrino portal (the composite SM singlet formed by the lepton doublet and the Higgs field). This interaction is such that it vanishes when the spin-3/2 field is on shell. In Sect. 2 below we give the model and basic constraints on the spin-3/2 field.

In Sect. 3 we study collider signatures of the spin-3/2 field. We study there $\nu_L h \rightarrow \nu_L h$ and $e^- e^+ \rightarrow W^+ W^-$ scatterings in detail. We give analytical computations and numerical predictions. We propose there a neutrino–Higgs collider and emphasize the importance of the linear collider in probing the spin-3/2 field.

In Sect. 4 we turn to loop effects of the spin-3/2 field. We find that the spin-3/2 field adds logarithmic and quartic UV-sensitivities atop the logarithmic and quadratic ones in the SM. We convert power-law UV-dependent terms into curvature terms as a result of the incorporation of gravity into the SM. Here we use the results of [5–7], which show that gravity can be incorporated into the SM properly and naturally (i) if the requisite curved geometry is structured by interpreting the UV cutoff as a constant value assigned to the spacetime curvature, and (ii) if the SM is extended by a secluded new physics (NP) that does not have to interact with the SM. This mechanism eliminates the big hierarchy problem by metamorphosing the quadratic UV part of the Higgs boson mass turns into Higgs–curvature coupling.

In Sect. 5 we discuss the possibility of Higgs inflation via the large Higgs non-minimal coupling induced by the spin-3/2 field. We find that Higgs inflation is possible in a wide range of parameters provided that the secluded NP sector is crowded enough.

In Sect. 6 we discuss the DM. We show therein that the spin-3/2 field is a viable DM candidate. We also show that the singlet fields in the NP can form a non-interacting DM component.

In Sect. 7 we conclude. There, we give a brief list of problems that can be studied as furthering of the material presented in this work.

2 A light spin-3/2 field

Introduced for the first time by Rarita and Schwinger [2], ψ_μ propagates with

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$$S^{\alpha\beta}(p) = \frac{i}{\not{p} - M}\Pi^{\alpha\beta}(p), \quad (1)$$

carrying one spin-3/2 and two spin-1/2 components through the projector [3,4],

$$\Pi^{\alpha\beta} = -\eta^{\alpha\beta} + \frac{\gamma^\alpha\gamma^\beta}{3} + \frac{\left(\gamma^\alpha p^\beta - \gamma^\beta p^\alpha\right)}{3M} + \frac{2p^\alpha p^\beta}{3M^2}, \quad (2)$$

which exhibits both spinor and vector characteristics. It is necessary to impose [3,4]

$$p^\mu \psi_\mu(p)\rfloor_{p^2=M^2} = 0 \quad (3)$$

and

$$\gamma^\mu \psi_\mu(p)\rfloor_{p^2=M^2} = 0, \quad (4)$$

to eliminate the two spin-1/2 components to make ψ_μ satisfy the Dirac equation

$$\left(\not{p} - M\right)\psi_\mu = 0 \quad (5)$$

as expected of an on-shell fermion. The constraints (3) and (4) imply that $p^\mu\psi_\mu(p)$ and $\gamma^\mu\psi_\mu(p)$ both vanish on the physical shell $p^2 = M^2$. The latter is illustrated in Fig. 1 taking ψ_μ on shell.

Characteristic of singlet fermions, the ψ_μ, at the renormalizable level, makes contact with the SM via

$$\mathcal{L}_{3/2}^{(int)} = c_{3/2}^i \overline{L^i} H \gamma^\mu \psi_\mu + \text{h.c.} \quad (6)$$

in which

$$L^i = \begin{pmatrix} \nu_{\ell L} \\ \ell_L \end{pmatrix}_i \quad (7)$$

is the lepton doublet ($i = 1, 2, 3$) and

$$H = \frac{1}{\sqrt{2}}\begin{pmatrix} v + h + i\varphi^0 \\ \sqrt{2}\varphi^- \end{pmatrix} \quad (8)$$

is the Higgs doublet with vacuum expectation value $v \approx 246$ GeV, Higgs boson h, and Goldstone bosons φ^-, φ^0 and φ^+ (forming the longitudinal components of W^-, Z and W^+ bosons, respectively).

In general, neutrinos are sensitive probes of singlet fermions. They can get masses through, for instance, the Yukawa interaction (6), which leads to the Majorana mass matrix

$$(m_\nu)_{3/2}^{ij} \propto c_{3/2}^i \frac{v^2}{M} c_{3/2}^{\star j} \quad (9)$$

after integrating out ψ_μ. This mass matrix, however, cannot lead to the experimentally known neutrino mixings [8–10]. This means that flavor structures necessitate additional singlet fermions. Of such a type are the right-handed neutrinos ν_R^k of mass M_k ($k = 1, 2, 3, \ldots$), which interact with the SM through

$$\mathcal{L}_R^{(int)} = c_R^{ik} \bar{L}^i H \nu_R^k + \text{h.c.} \quad (10)$$

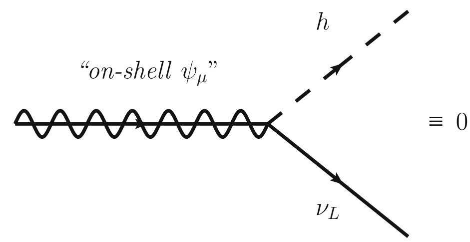

Fig. 1 ψ_μ–h–ν_L coupling with vertex factor $ic_{3/2}\gamma^\mu$. Scatterings in which ψ_μ is on shell must all be forbidden since $c_{3/2}\gamma^\mu\psi_\mu$ vanishes on mass shell by the constraint (4). This ensures stability of ψ_μ against decays and all sort of co-annihilations

to generate the neutrino Majorana masses

$$(m_\nu)_R^{ij} \propto c_R^{ik} \frac{v^2}{M_k} c_R^{\star kj} \quad (11)$$

of more general flavor structure. This mass matrix must have enough degrees of freedom to fit to the data [8–10].

Here we make a pivotal assumption. We assume that ψ_μ and ν_R^k can weigh as low as a TeV, and that $c_{3/2}^i$ and some of c_R^{ik} can be $\mathcal{O}(1)$. We, however, require that contributions to neutrino masses from ψ_μ and ν_R add up to reproduce the experimental result

$$(m_\nu)_{3/2}^{ij} + (m_\nu)_R^{ij} \approx (m_\nu)_{\text{exp}}^{ij} \quad (12)$$

via cancellations among different terms. We therefore take

$$c_{3/2} \lesssim \mathcal{O}(1), \quad M \gtrsim \text{TeV} \quad (13)$$

and investigate the physics of ψ_μ. This cancellation requirement does not have to cause any excessive fine-tuning simply because ψ_μ and ν_R^k can have appropriate symmetries that correlate their couplings. One possible symmetry would be rotation of $\gamma^\mu\psi_\mu$ and ν_R^k into each other. We defer study of possible symmetries to another work, which is in progress [11]. The right-handed sector, which can involve many ν_R^k fields, is interesting by itself but hereon we focus on ψ_μ and take, for simplicity, $c_{3/2}^i$ real and family-universal ($c_{3/2}^i = c_{3/2}$ for $\forall\, i$).

3 Spin-3/2 field at colliders

It is only when it is off-shell that ψ_μ can reveal itself through the interaction (6). This means that its effects are restricted to modifications in scattering rates of the SM particles. To this end, as follows from (6), it participates in

1. $\nu_L h \to \nu_L h$ (and also $\nu_L\nu_L \to hh$),
2. $e^+e^- \to W_L^+ W_L^-$ (and also $\nu_L\nu_L \to Z_L Z_L$),

at the tree level. They are analyzed below in detail.

3.1 $\nu_L h \to \nu_L h$ Scattering

Shown in Fig. 2 are the two box diagrams which enable $\nu_L h \to \nu_L h$ scattering in the SM. Added to this loop-suppressed SM piece is the ψ_μ piece depicted in Fig. 3. The two contributions add up to give the cross section

$$\frac{\mathrm{d}\sigma(\nu_L h \to \nu_L h)}{\mathrm{d}t} = \frac{1}{16\pi}\frac{\mathcal{T}_{\nu h}(s,t)}{(s-m_h^2)^2}, \tag{14}$$

in which the squared matrix element

$$\begin{aligned}\mathcal{T}_{\nu h}(s,t) = 9\left(\frac{c_{3/2}}{3M}\right)^4\left(\left(s-m_h^2\right)^2+st\right)\\ -16\left(\frac{c_{3/2}}{3M}\right)^2\left(2\left(s-m_h^2\right)^2+\left(2s-m_h^2\right)t\right)\\ \times\mathbb{L}+2\left(s-m_h^2\right)\left(s+t-m_h^2\right)\mathbb{L}^2\end{aligned} \tag{15}$$

involves the loop factor

$$\mathbb{L} = \frac{(g_W^2+g_Y^2)^2 M_Z^2 m_h^2 I(M_Z)}{192\pi^2} + \frac{g_W^4 M_W^2 m_h^2 I(M_W)}{96\pi^2} \tag{16}$$

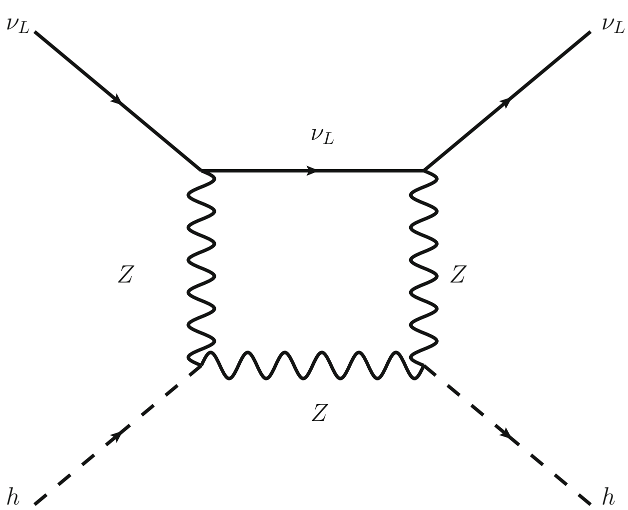

Fig. 2 The ν–Z box mediating the $\nu_L h \to \nu_L h$ scattering in the SM. The e–W box is not shown

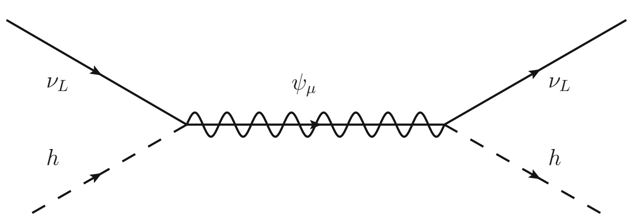

Fig. 3 $\nu_L h \to \nu_L h$ scattering with ψ_μ mediation. No resonance can occur at $\sqrt{s} = M$ because ψ_{DM} cannot come to mass shell

in which g_W (g_Y) is the isospin (hypercharge) gauge coupling, and

$$\begin{aligned}I(\mu) &= \int_0^1 \mathrm{d}x \int_0^{1-x}\mathrm{d}y\int_0^{1-x-y}\\ &\times \mathrm{d}z\Big((s-m_h^2)(x+y+z-1)y\\ &-txz+m_h^2 y(y-1)+\mu^2(x+y+z)\Big)^{-2}\end{aligned} \tag{17}$$

is the box function. In Fig. 4, we plot the total cross section $\sigma(\nu_L h \to \nu_L h)$ as a function of the neutrino–Higgs center-of-mass energy for different M values. The first important thing about the plot is that there is no resonance formation around $\sqrt{s} = M$. This confirms the fact that ψ_μ, under the constraint (4), cannot come to physical shell with the couplings in (6). In consequence, the main search strategy for ψ_μ is to look for deviations from the SM rates rather than resonance shapes. The second important thing about the plot is that, in general, as revealed by (19), the larger M, the smaller the ψ_μ contribution. The cross section starts around 10^{-7} pb, and falls rapidly with $\sqrt{s}$. (The SM piece, as a loop effect, is too tiny to be observable: $\sigma(\nu_L h \to \nu_L h) \lesssim 10^{-17}$ pb.) It is necessary to have some $10^4/fb$ integrated luminosity (100 times the target luminosity at the LHC) to observe a few events in a year. This means that $\nu_L\nu_L \to hh$ scattering can probe ψ_μ only at high luminosity but with a completely new scattering scheme.

Figure 4 shows that neutrino–Higgs scattering can be a promising channel to probe ψ_μ (at high-luminosity, high-energy machines). The requisite experimental setup would involve crossing of Higgs factories with accelerator neutrinos. The setup, schematically depicted in Fig. 5, can be viewed as incorporating future Higgs (CEPC [12], FCC-ee [13,14] and ILC [15–17]) and neutrino [18–20] factories. If

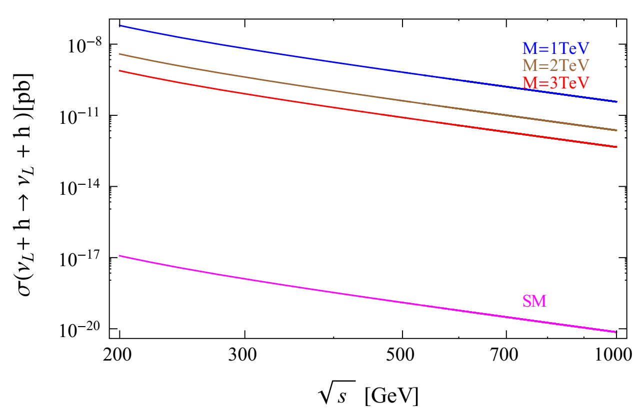

Fig. 4 The total cross section for $\nu_L h \to \nu_L h$ scattering as a function of the neutrino–Higgs center-of-mass energy $\sqrt{s}$ for $M = 1, 2$ and 3 TeV at $c_{3/2} = 1$. Cases with $c_{3/2} \neq 1$ can be reached via the rescaling $M \to M/c_{3/2}$

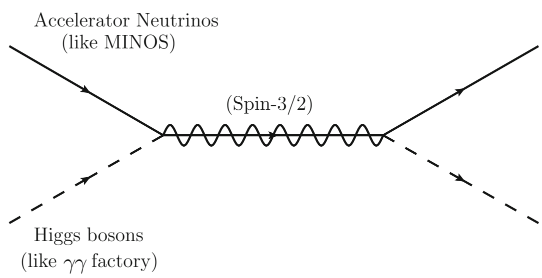

Fig. 5 Possible neutrino–Higgs collider to probe ψ_μ

ever realized, it could be a rather clean experiment with negligible SM background. This hypothetical "neutrino–Higgs collider", depicted in Fig. 5, must have, as suggested by Fig. 4, some $10^4/fb$ integrated luminosity to be able to probe a TeV-scale ψ_μ. In general, the need to have high luminosities is a disadvantage of this channel. (A feasibility study, technical design and possible realization of a "neutrino–Higgs collider" fall outside the scope of the present work.)

3.2 $e^+e^- \to W_L^+ W_L^-$ Scattering

It is clear that ψ_μ directly couples to the Goldstone bosons $\varphi^{+,-,0}$ via (6). The Goldstones, though eaten up by the W and Z bosons in acquiring their masses, reveal themselves at high energies. In fact, the Goldstone equivalence theorem [21–23] states that scatterings at energy E involving longitudinal $W_L^\pm$ bosons are equal to scatterings that involve $\varphi^\pm$ up to terms $\mathcal{O}(M_W^2/E^2)$. This theorem, with similar equivalence for the longitudinal Z boson, provides a different way of probing ψ_μ. In this regard, depicted in Fig. 6 is ψ_μ contribution to $e^+e^- \to W_L^+ W_L^-$ scattering in light of the Goldstone equivalence. The SM amplitude is given in [21–23]. The total differential cross section

$$\frac{d\sigma(e^+e^- \to W_L^+ W_L^-)}{dt} = \frac{1}{16\pi s^2}\mathcal{T}_{W_L W_L}(s,t) \tag{18}$$

involves the squared matrix element

$$\begin{aligned}\mathcal{T}_{W_L W_L}(s,t) = &\left(\frac{g_W^2}{s-M_Z^2}\left(-1+\frac{M_Z^2}{4M_W^2}+\frac{M_Z^2-M_W^2}{s}\right)\right.\\ &\left.+\frac{g_W^2}{s-4M_Z^2}\left(1+\frac{M_W^2}{t}\right)+\frac{c_{3/2}^2}{3M^2}\right)^2\\ &\times\left(-2sM_W^2-2(t-M_W^2)^2\right)\\ &+\frac{c_{3/2}^4 s}{18M^2}\left(4+\frac{t}{t-M^2}\right)^2.\end{aligned} \tag{19}$$

Plotted in Fig. 7 is $\sigma(e^+e^- \to W_L^+ W_L^-)$ as a function of the e^+e^- center-of-mass energy for different values of M. The cross section, which falls with $\sqrt{s}$ without exhibiting a resonance shape, is seen to be large enough to be measurable at the ILC [15–17]. In general, the larger M, the smaller the cross section but even $1/fb$ luminosity is sufficient for probing ψ_μ for a wide range of mass values.

Collider searches for ψ_μ, as illustrated by $\nu_L h \to \nu_L h$ and $e^-e^+ \to W^+W^-$ scatterings, can access spin-3/2 fields of several TeV mass. For instance, the ILC, depending on its precision, can confirm or exclude a ψ_μ of even 5 TeV mass with an integrated luminosity around $1/fb$. Depending on the possibility and feasibility of a neutrino–neutrino collider (mainly accelerator neutrinos), it may be possible to study also $\nu_L\nu_L \to hh$ and $\nu_L\nu_L \to Z_L Z_L$ scatterings, which are expected to have similar sensitivities to M.

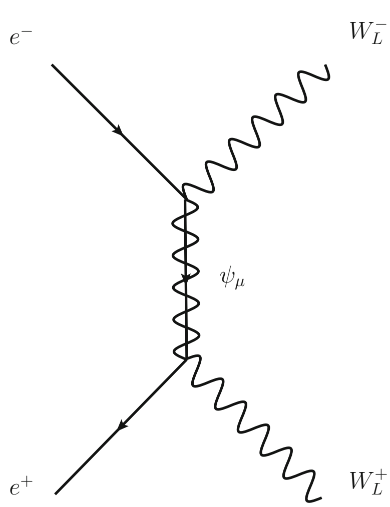

Fig. 6 The Feynman diagram for $e^+e^- \to W_L^+ W_L^-$ scattering. The $\nu_L\nu_L \to Z_L Z_L$ scattering has the same topology

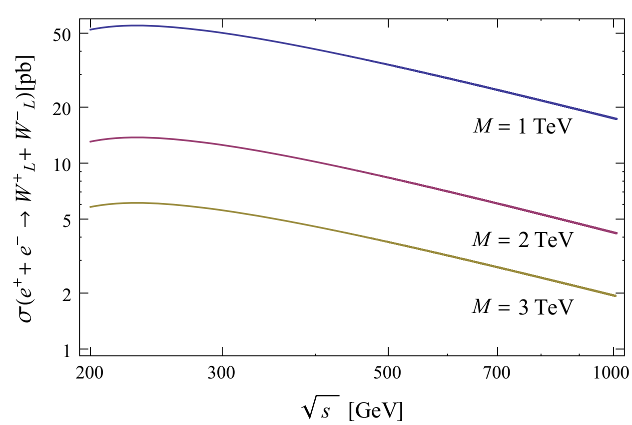

Fig. 7 The total cross section for $e^-e^+ \to W^+W^-$ scattering as a function of the electron–positron center-of-mass energy $\sqrt{s}$ for $M = 1, 2$ and 3 TeV at $c_{3/2} = 1$. Cases with $c_{3/2} \neq 1$ can be reached via the rescaling $M \to M/c_{3/2}$

4 Spin-3/2 field in loops

As an inherently off-shell field, ψ_μ is expected to reveal itself mainly in loops. Its only possible loop effect would be generation of neutrino masses but chirality forbids it. Despite the couplings in (6), therefore, neutrino masses do not get any contribution from the ψ_μ–h loop.

One other loop effect of ψ_μ would be radiative corrections to the Higgs boson mass. This is not forbidden by any symmetry. The relevant Feynman diagram is depicted in Fig. 8. It adds to the Higgs boson squared-mass a logarithmic piece,

$$\left(\delta m_h^2\right)_{log} = \frac{c_{3/2}^2}{12\pi^2} M^2 \log G_F M^2, \tag{20}$$

relative to the logarithmic piece $\log G_F \Lambda^2$ in the SM and a quartic piece

$$\left(\delta m_h^2\right)_4 = \frac{c_{3/2}^2}{48\pi^2} \frac{\Lambda^4}{M^2}, \tag{21}$$

which have the potential to override the experimental result [1] depending on how large the UV cutoff Λ is compared to the Fermi scale $G_F^{-1/2} = 293$ GeV.

The logarithmic contribution in (20), which originates from the $\eta^{\alpha\beta}$ part of (2), gives rise to the little hierarchy problem in that the larger M, the stronger the destabilization of the SM Higgs sector. Leaving aside the possibility of cancellations with similar contributions from the right-handed neutrinos ν_R^k in (10), the little hierarchy problem can be prevented if M (more precisely $M/c_{3/2}$) lies in the TeV domain.

The quartic contribution in (21), which originates from the longitudinal $p^\alpha p^\beta$ term in (2), gives cause to the notorious big hierarchy problem in that the larger Λ is, the larger the destabilization of the SM Higgs sector. This power-law UV sensitivity exists already in the SM

$$\left(\delta m_h^2\right)_2 = \frac{3\Lambda^2}{16\pi^2 |\langle H \rangle|^2} \left(m_h^2 + 2M_W^2 + M_Z^2 - 4m_t^2\right) \tag{22}$$

at the quadratic level [24] and violates the LHC bounds unless $\Lambda \lesssim 550$ GeV. This bound obviously contradicts with the LHC experiments since the latter continue to confirm the

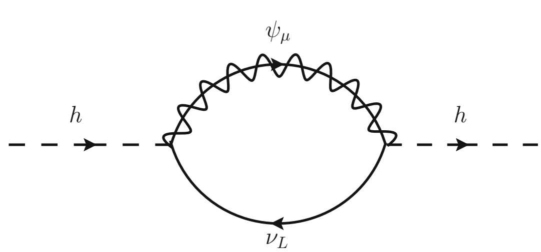

Fig. 8 The $\psi_\mu - \nu_L$ loop that generates the logarithmic correction in (20) and the quartic correction in (21)

SM at multi TeV energies. This experimental fact makes it obligatory to find a natural UV completion to the SM.

One possibility is to require $\left(\delta m_h^2\right)_4$ to cancel out $\left(\delta m_h^2\right)_2$. This requirement involves a severe fine-tuning (as with a scalar field [31–34], Stueckelberg vector [35] and spacetime curvature [36,37]) and cannot form a viable stabilization mechanism.

Another possibility would be to switch, for instance, to a dimensional regularization scheme, wherein the quartic and quadratic UV-dependencies are known to disappear. This, however, is not a solution. The reason is that the SM, as a quantum field theory of the strong and electroweak interactions, needs gravity to be incorporated as the fourth known force. And the fundamental scale of gravity, M_{Pl}, inevitably sets an non-eliminable physical UV cutoff (rendering Λ physical). This cutoff forces quantum field theories to exist in between physical UV and IR scales. The SM plus ψ_μ (plus right-handed neutrinos), for instance, ranges from $G_F^{-1/2}$ at the IR up to Λ at the UV such that both scales are physical (not to be confused with the formal momentum cutoffs employed in the cutoff regularization).

To stabilize the SM, it is necessary to metamorphose the destabilizing UV effects. This necessitates a physical agent. The most obvious candidate is gravity. That is to say, the UV-naturalness problems can be a clue to how quantized matter must gravitate. Indeed, quantized matter in classical curved geometry suffers from inconsistencies. The situation can be improved by considering long-wavelength matter by integrating out high-frequency modes. This means that the theory to be carried into curved geometry for incorporating gravity is not the full action but the effective action (see the discussions in [5–7]). Thus, starting with the SM effective action in flat spacetime with well-known logarithmic, quartic and quadratic UV-sensitivities, gravity can be incorporated in a way that ensures UV-naturalness. More precisely, gravity gets incorporated properly and naturally (i) if the requisite curved geometry is structured by interpreting Λ^2 as a constant value assigned to the spacetime curvature, and (ii) if the SM is extended by new physics (NP) that does not have to interact with the SM. The ψ_μ can well be an NP field. Incorporating gravity by identifying $\Lambda^2 g_{\mu\nu}$ with the Ricci curvature $R_{\mu\nu}(g)$, fundamental scale of gravity gets generated as

$$M_{Pl}^2 \approx \frac{(n_b - n_f)}{2(8\pi)^2} \Lambda^2 \tag{23}$$

where n_b (n_f) are the total number of bosons (fermions) in the SM plus the NP. The ψ_μ increases n_f by 4, right-handed neutrinos by 2. There are various other fields in the NP, which contribute to n_b and n_f to ensure $\Lambda \lesssim M_{Pl}$. Excepting ψ_μ, they do not need to interact with the SM fields. Induction of M_{Pl} ensures that the quadratic UV-contributions to vacuum energy are canalized not to the cosmological constant but to the gravitational constant (see [38,39], arriving at this result in a different context). This suppresses the cosmological constant down to the neutrino mass scale.

The quartic UV-contributions in (21) and the quadratic contributions in (22) (suppressing contributions from the right-handed neutrinos ν_R^k) change their roles with the inclusion of gravity. Indeed, corrections to the Higgs mass term $\left[(\delta m_h^2)_4 + (\delta m_h^2)_2\right] H^\dagger H$ turn into

$$\left[\frac{3\left(m_h^2+2M_W^2+M_Z^2-4m_t^2\right)}{(8\pi)^2\,|\langle H\rangle|^2}+\frac{c_{3/2}^2}{12(n_b-n_f)}\frac{M_{Pl}^2}{M^2}\right]R H^\dagger H, \tag{24}$$

which is nothing but the direct coupling of the Higgs field to the scalar curvature R. This Higgs–curvature coupling is perfectly natural; it has no potential to de-stabilize the Higgs sector. Incorporation of gravity as in [5–7] leads, therefore, to UV-naturalization of the SM with a nontrivial NP sector containing ψ_μ as its interacting member.

5 Spin-3/2 field as enabler of higgs inflation

The non-minimal Higgs–curvature coupling in (24) reminds one at once of the possibility of Higgs inflation. Indeed, the Higgs field has been shown in [40,41] to lead to correct inflationary expansion provided that

$$\frac{c_{3/2}^2}{12(n_b-n_f)}\frac{M_{Pl}^2}{M^2} \approx 1.7\times 10^4 \tag{25}$$

after dropping the small SM contribution in (24). This relation puts constraints on M and Λ, depending on how crowded the NP is.

For a Planckian UV cutoff $\Lambda \approx M_{Pl}$, the Planck scale in (23) requires $n_b - n_f \approx 1300$, and this leads to $M/c_{3/2} \approx 6.3\times 10^{13}$ GeV. This heavy ψ_μ, weighing not far from the see-saw and axion scales, acts as an enabler of Higgs inflation. (Of course, all this makes sense if the ψ_μ contribution in (20) is neutralized by similar contributions from the right-handed neutrinos ν_R^k to alleviate the little hierarchy problem.)

For an intermediate UV cutoff $\Lambda \ll M_{Pl}$, $n_b - n_f$ can be large enough to bring M down to lower scales. In fact, M gets lowered to $M \sim$ TeV for $n_b - n_f \simeq 10^{24}$, and this sets the UV cutoff $\Lambda \sim 3$ TeV. This highly crowded NP illustrates how small M and Λ can be. Less crowded NP sectors lead to intermediate-scale M and Λ.

It follows therefore that it is possible to realize Higgs inflation through the Higgs–curvature coupling (corresponding to quartic UV-dependence the ψ_μ induces on the Higgs mass). It turns out that Higgs inflation is decided by how heavy ψ_μ is and how crowded the NP is. It is interesting that the ψ_μ hidden in the SM spectrum enables successful Higgs inflation if gravity is incorporated into the SM as in [5–7].

6 Spin-3/2 field as dark matter

Dark matter (DM), forming one-fourth of the matter in the Universe, must be electrically neutral and long-lived. The negative searches [42–44] so far have added one more feature: The DM must have exceedingly suppressed interactions with the SM matter. It is not hard to see that the spin-3/2 fermion ψ_μ possesses all these properties. Indeed, the constraint (4) ensures that scattering processes in which ψ_μ is on its mass shell must all be forbidden simply because its interaction in (6) involves the vertex factor $c_{3/2}\gamma^\mu$. This means that decays of ψ_μ as in Fig. 1 as well as its co-annihilations with the self- and other SM fields are all forbidden. Its density therefore does not change with time, and the observed DM relic density [45] must be its primordial density, which is determined by the short-distance physics the ψ_μ descends from. It is not possible to calculate the relic density without knowing the short-distance physics. Its mass and couplings, on the other hand, can be probed via the known SM-scatterings as studied in Sect. 3 above. In consequence, the ψ_μ, as an inherently off-shell fermion hidden in the SM spectrum, possesses all the features required of a DM candidate.

Of course, the ψ_μ is not the only DM candidate in the setup. The crowded NP sector, needed to incorporate gravity in a way solving the hierarchy problem (see Sect. 4 above), involves various fields which do not interact with the SM matter. They are viable candidates for non-interacting DM as well as dark energy (see the detailed analysis in [7]). The non-interacting NP fields can therefore contribute to the total DM distribution in the Universe. It will, of course, not be possible to search for them directly or indirectly. In fact, they do not have to come to equilibrium with the SM matter.

Interestingly, both ψ_μ and the secluded fields in the NP act as extra fields hidden in the SM spectrum. Unlike the ψ_μ, which reveal itself virtually, the NP singlets remain completely intact. The main implication is that, in DM phenomenology, one must keep in mind that there can exist an unobservable, undetectable component of the DM [7].

7 Conclusion and outlook

In this work we have studied a massive spin-3/2 particle ψ_μ obeying the constraint (4) and interacting with the SM via (6). It hides in the SM spectrum as an inherently off-shell field. We first discussed its collider signatures by studying $\nu_L h \rightarrow \nu_L h$ and $e^- e^+ \rightarrow W^+ W^-$ scatterings in detail in Sect. 3. Following this, we turned to its loop effects and determined how it contributes to the big and little hierarchy problems in the SM. Resolving the former by appropriately incorporating gravity, we show that the Higgs field can inflate the Universe. Finally, we show that ψ_μ is a viable DM candidate, which can be indirectly probed via the scattering processes we have analyzed.

The material presented in this work can be extended in various ways. A partial list would include:

- Determining under what conditions right-handed neutrinos can lift the constraints on ψ_μ from the neutrino masses.
- Improving the analyses of $\nu_L h \to \nu_L h$ and $e^- e^+ \to W^+ W^-$ scatterings by including loop contributions.
- Simulating $e^- e^+ \to W^+ W^-$ at the ILC by taking into account planned detector acceptances and collider energies.
- Performing a feasibility study of the proposed neutrino–Higgs collider associated with $\nu_L h \to \nu_L h$ scattering.
- Exploring UV-naturalness by including right-handed neutrinos, and determining under what conditions the little hierarchy problem is softened.
- Including effects of the right-handed neutrinos into Higgs inflation, and determining appropriate parameter space.
- Giving an in-depth analysis of the dark matter and dark energy by taking into account the spin-3/2 field, right-handed neutrinos and the secluded NP fields.
- Studying constraints on the masses of NP fields from nucleosynthesis and other processes in the early Universe.

We will continue to study the spin-3/2 hidden field starting with some of these points.

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