The Hubble tension and a renormalizable model of gauged neutrino self-interactions

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We present a simple extension of the Standard Model that leads to renormalizable long-range vector-mediated neutrino self-interactions. This model can resolve the Hubble tension by delaying the onset of neutrino free-streaming during recombination, without conflicting with other measurements. The extended gauge, scalar and neutrino sectors lead to observable signatures, including invisible Higgs and Z decays, thereby relating the Hubble tension to precision measurements at the LHC and future colliders. The model has a new neutrino-philes gauge boson with $m_{Z'} \sim O(10\text{ eV})$ and charged Higgses at a few 100 GeV. It requires hidden neutrinos with active-hidden mixing angles larger than $5 \times 10^{-8}$ and masses in the range $1 \div 300\text{ eV}$, which could also play a role for short baseline neutrino oscillation anomalies.

Introduction.—There is convincing evidence that neutrinos played a substantial role during the epoch of big bang nucleosynthesis (BBN) at $T \sim \text{MeV}$, closely monitored by early element abundances. The lowest temperature scale indirectly probed for neutrinos is $T \sim \text{eV}$, where observations of the cosmic microwave background (CMB) fit well to a history of our universe that does not only comply with the cosmological standard model (ΛCDM), but also with the expectation of the Standard Model of particle physics (SM), including exactly three generations of neutrinos.

However, evidence is accumulating not only for a discrepancy between local measurements of today’s Hubble rate $H_0$ [1–5] and therelike global determinations based on ΛCDM together with CMB [6], baryonic acoustic oscillations (BAO) and large scale structure (LSS) datasets [7–16], but also for an increasing tension in other parameters, see e.g. [17, 18]. The ultimate resolution of those discrepancies might require a modification of ΛCDM, preferentially, perhaps, shortly before the era of recombination [19, 20]. Too many new physics (NP) scenarios have been discussed to review all of them, see [20–23] and references therein. Naturally, any consistent modification of ΛCDM must be in compliance with a consistent modification of the SM.

The positive correlation of $H_0$ and $N_{\text{eff}}$ with the amplitude of the matter power spectrum $σ_8$, as observed in CMB data [6], prohibits a resolution of the $H_0$ tension simply by increasing $N_{\text{eff}}$ alone (LSS prefers low $σ_8$). However, a delay in the onset of neutrino free streaming during recombination could achieve both: breaking the positive correlation of $H_0$ and $σ_8$, while solving the Hubble tension at the cost of increasing $\Delta N_{\text{eff}}$ during recombination [24–29]. Taking into account an effective four-neutrino interaction $G_{\text{eff}}^{4\nu}(\bar{\nu} \nu)(\bar{\nu} \nu)$ with strength $G_{\text{eff}}^{4\nu}$ a good, bi-modal fit to CMB data is obtained with [28, 29]

$$G_{\text{eff}}^{4\nu} \equiv \frac{g_{\text{eff}}^2}{m_{Z'}^2} \approx \begin{cases} (5\text{ MeV})^{-2} & \text{(SI)}, \\
(100\text{ MeV})^{-2} & \text{(WI)} \end{cases}$$

The weakly interacting mode (WI) should be interpreted as an upper limit on $G_{\text{eff}}^{4\nu}$ such that the fit of cosmological parameters stays close to ΛCDM [24, 25], which then of course does not resolve above tensions. Therefore, we will focus on the strongly interacting mode (SI), which considerably alters cosmology to resolve the tensions in $H_0$ and $σ_8$ while being consistent with local astronomical observations [28, 29].

While decoupled heavy new physics (NP) certainly is suited to generate $G_{\text{eff}}^{4\nu}$, we stress that it is not precluded that also a light mediator with interaction strength similar to the SI mode might resolve the tensions [30, 31]. But of course, light mediators strongly interacting with neutrinos are highly constrained by the bound on $\Delta N_{\text{eff}}$ during BBN, see e.g. [32]. However, while one may feel that it is just a relatively short time between BBN and recombination, we recall that it is still six orders of magnitude in temperature. This certainly is enough to establish a mass scale, say after a phase transition, and subsequently integrate it out to obtain a decoupling behavior of neutrinos during CMB closely resembling (1). In this way, neutrinos recouple by the new interactions only after BBN, and fall out of equilibrium shortly before or during recombination.

In this letter we provide what we think is the simplest renormalizable and phenomenologically viable extension of the SM that leads to vector mediated four-neutrino interactions of above strength. We first outline the parameter space suitable to address the Hubble tension. Subsequently we fully flesh out our model, discuss constraints on the parameter space and means to test the model.

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Relevant parameter region.—The effective four-neutrino interaction strength in our model is
\[
G_{4\nu}^{\text{eff}} \equiv \frac{g_4^{\text{eff}}}{m_{Z'}^2} \equiv \frac{g_X^2 v_m^2}{m_{Z'}^2},
\]
where \( g_X \) is the gauge coupling of a new U(1)\(_X\) symmetry, \( \varepsilon_m \ll 1 \) a mixing between active and hidden (U(1)\(_X\) charged) neutrinos, and \( m_{Z'} \) the mass of the new gauge boson after U(1)\(_X\) breaking. Equating the resulting thermally averaged interaction rate with the Hubble rate \( H \sim T^2/M_{\text{Pl}} \) (for radiation dominated universe – results only change mildly if we switch to matter domination for \( T \lesssim 0.8 \text{ eV} \)) one finds
\[
(G_{4\nu}^{\text{eff}})^2 T^5 \approx T^2/M_{\text{Pl}}.
\]
Using (1), the decoupling temperature is obtained as \( T_{\text{dec}} \approx 0.5 \text{ eV} \) confirming this interaction freezes out around recombination.

On the other hand, for a range of temperatures \( T \gg m_{Z'} \), while \( \varepsilon_m \) is relevant, the new gauge boson will be effectively massless giving rise to an induced long-range four-neutrino interaction with thermally averaged rate \( \Gamma \sim g_4^{\text{eff}} T \). Requiring this interaction not to thermalize with neutrinos prior to BBN, but before recombination, results in a narrow parameter window
\[
2 \times 10^{-7} \lesssim g_X^2 \varepsilon_m^2 \lesssim 5 \times 10^{-6}.
\]
Knowing \( g_{\text{eff}} \) and \( G_{4\nu}^{\text{eff}} \) we can compute \( m_{Z'} \) to find
\[
1 \lesssim m_{Z'} \lesssim 25 \text{ eV} \quad (\text{SI}).
\]
Furthermore, parametrizing \( m_{Z'} = g_X \bar{v} \) we can also constrain the size of the effective U(1)\(_X\) breaking vacuum expectation value (VEV) \( \bar{v} \) to
\[
\bar{v} := \frac{m_{Z'}}{g_X} \approx \varepsilon_m^2 \times 5 \text{ MeV} \quad (\text{SI}).
\]
We illustrate this particular re- and decoupling behavior in Fig. 1, computed exactly within our model for representative parameter points. Note that (6) implies a hierarchy between the relevant scales in the model of
\[
\xi := \bar{v}/v_h \approx \varepsilon_m^2 \times 2 \times 10^{-5} \quad (\text{SI}),
\]
where \( v_h = 246 \text{ GeV} \) is the SM Higgs VEV.

The Model.— Next to the new U(1)\(_X\) gauge symmetry we introduce a pair of SM-neutral chiral fermions \( N_{1,2} \) and two new scalars \( \Phi \) and \( S \) with charges as shown in Tab. I [102]. New interaction terms for SM leptons are given by
\[
\mathcal{L}_{\text{new}} = -y \bar{L} \Phi N_1 - M N_1 N_2 + \text{h.c.},
\]
where \( \Phi := i \sigma_2 \Phi^* \), \( y \) is a dimensionless Yukawa coupling, and \( M \) has mass-dimension one. For brevity, we stick to the one generation case here, while extensions to three

\[
\begin{array}{c|c|c|c|c}
\text{Field} & \Phi & N_1 & N_2 & S \\
\hline
\text{Lorentz} & S & \text{RH} & \text{RH} & S \\
\text{SU}(2)_L \times U(1)_Y & (2, -\frac{1}{2}) & 0 & 0 & 0 \\
U(1)_{X} & +1 & +1 & -1 & +1 & 0 \\
U(1)_{L} & 0 & +1 & -1 & 0 & 0
\end{array}
\]

FIG. 1: Thermally averaged four-neutrino interaction rate relative to the Hubble rate as a function of Temperature for \( m_{Z'} = 25 \text{ eV} \) and two different values of the \( Z' \) width.

TABLE I: New fields and their charges under Lorentz, SM gauge, new U(1)\(_X\) gauge symmetry as well as under global Lepton number (S=Scalar, RH=right-handed Weyl fermion, V=vector).

generations of SM leptons or multiple generations of hidden fermions are straightforward, and considered below. The most general scalar potential consistent with all symmetries is
\[
V = V_H + V_\Phi + V_S + V_{H\Phi} + V_H S + V_\Phi S + V_3,
\]
with
\[
\begin{align*}
V_2 &:= \mu^2 \Sigma^\dagger \Sigma + \lambda \Sigma (\Sigma^\dagger \Sigma)^2 \quad (\Sigma = H, \Phi, S), \\
V_{H\Phi} &:= \lambda_3 (H^\dagger H) (\Phi^\dagger \Phi) + \lambda_4 (H^\dagger \Phi) (\Phi^\dagger H), \\
V_{DS} &:= \lambda_{DS} (D^\dagger D) (S^\dagger S) \quad (D = H, \Phi), \\
V_3 &:= -\sqrt{2} \mu (H^\dagger \Phi) S^* + \text{h.c.}.
\end{align*}
\]
We decompose the scalars as
\[
H = \begin{pmatrix} h^+ \\ \frac{1}{\sqrt{2}} (h + i a_h) \end{pmatrix}, \quad \Phi = \begin{pmatrix} \phi^+ \\ \frac{1}{\sqrt{2}} (\phi + i a_\phi) \end{pmatrix},
\]
and \( S = \frac{1}{\sqrt{2}} (s + i a_s) \).

We choose a parameter region such that all neutral scalars obtain VEVs \( v_\sigma := \langle \sigma \rangle \) for \( \sigma = h, \phi, s \), and assume CP conservation in the scalar sector. \( v_h \) spontaneously
breaks EW symmetry, \( v_s \) breaks the new U(1)\(_X\), while \( v_\phi \) breaks both. Fixing the Hubble tension requires the hierarchy \( v_h \gg v_s, v_\phi \), cf. (7), and we will expand all of our expressions to leading order in that hierarchy.

The photon is exactly the same massless combination of EW bosons as in the SM, mixed by the electroweak angle \( \gamma \) by \( m_W/m_Z \) \([103]\). By contrast, the very SM-like Z boson contains a miniscule admixture of the new gauge boson \( X \),

\[
Z_\mu = c_X (c_W W_\mu^3 - s_W B_\mu) + s_X X_\mu, \tag{16}
\]

with an angle \([104]\)

\[
s_X \approx -2 c_W \frac{gX}{g_2} \left( \frac{v_\phi}{v_h} \right)^2 \ll 1 \quad \text{and} \quad c_X \approx 1. \tag{17}
\]

The masses of the physical neutral gauge bosons up to \( O(\xi^2) \) are

\[
m_Z \approx \frac{g_2 v_h}{2 c_W} \quad \text{and} \quad m_{Z'} \approx g_X v := g_X \sqrt{v_\phi^2 + v_s^2}. \tag{18}
\]

Taking into account \( v_\phi \), the neutrino mass matrix in the gauge basis \((\nu, \nu_1, \nu_2)\) is given by

\[
M_\nu = \begin{pmatrix} 0 & -y v_\phi/\sqrt{2} & 0 \\ -y v_\phi/\sqrt{2} & 0 & M \\ 0 & M & 0 \end{pmatrix}. \tag{19}
\]

Upon 13–rotating by an angle \( \varepsilon_m \) with

\[
\tan \varepsilon_m := (y v_\phi)/(\sqrt{2} M), \tag{20}
\]

this matrix has an exact zero eigenvalue, corresponding to approximately massless active neutrinos, and a Dirac neutrino \( N \) with mass \( M_N := \sqrt{M^2 + y^2 v_\phi^2}/2 \). The massless active neutrinos mix with \( \nu_2 \) proportional to \( s_{\varepsilon_m} \) generating the coupling (2). Together with

\[
\tan \gamma := v_\phi/v_s, \tag{21}
\]

one can show that

\[
M = (y/\sqrt{2}) \varepsilon_m s_\gamma (G^{\nu}_{\text{eff}})^{-1/2} \ll 5 \text{ MeV}. \tag{22}
\]

Owing to constraints discussed below the parameter range one should have in mind is

\[
2 \times 10^{-5} \lesssim y \lesssim 6 \times 10^{-3}, \quad \varepsilon_m \lesssim 0.05 \quad \text{and} \quad s_\gamma \lesssim 0.2.
\]

The mass generation for active neutrinos \( m_\nu \ll v_\phi \) is merely a small perturbation to this setting. In particular, our mechanism is compatible with an effective Majorana mass in \([M_\nu]_{111}\), and, therefore, with any type of mass generation mechanism that gives rise to the Weinberg operator \([33]\). Another minimal possibility in the present model would be to populate \([M_\nu]_{333}\) like in the inverse seesaw mechanism \([34–36]\). Also Dirac masses are possible but require additional fermions. Ultimately, any of the commonly considered neutrino mass generation mechanisms is compatible with our model.

**Phenomenology.**—The scalar sector of the model corresponds to a 2HDM+scalar singlet. However, both of the new scalars are charged under the hidden-neutrino-specific U(1)\(_X\) which considerably alters phenomenology with respect to earlier works \([37–42]\). The masses of the physical scalars, to leading order in \( \xi \equiv \bar{v}/v_h \), are given by \([105]\)

\[
m^2_H = 2 \lambda_H v_h^2, \quad m^2_\Phi = \frac{2}{\xi^2} v_h \mu, \quad m^2_{h_s} \approx \frac{\lambda^2}{2} m^2_{h_s} + O(\gamma^2/v_h^2). \tag{23}
\]

We diagonalize the neutral scalar mass matrix by three orthogonal rotations \( O = R(\theta_13) R(\eta_{12}) R(\theta_{23}) \), such that

\[
O^T M^2_{h,s,n} = \text{Diag}(m^2_{h,s}, m^2_{h_s}, m^2_\Phi) \tag{24}
\]

The mixing angles, to leading order in \( \xi \), are given by

\[
s_{12} \equiv s_{\Phi} = s_\gamma, \quad s_{13} \equiv s_{H S} = \xi \frac{\mu}{2 v_h} \tag{27}
\]

\[
s_{23} \equiv s_{H} = \xi s_{\gamma} \frac{\mu_4}{2 \lambda_H} + \frac{1}{\lambda_H} v_h \left( \frac{\lambda_\mu S_\nu - \mu S_\nu}{2} \right), \tag{28}
\]

where we use \( \lambda_{34} := \lambda_3 + \lambda_4 \) and

\[
p := \lambda_{34} v_h \approx \mu S_\nu - \mu S_\nu, \quad q := \lambda H S v H c_\gamma - \mu S_\nu. \tag{29}
\]

For the parameter region envisaged to resolve the Hubble tension, there are two new light bosonic fields: next to \( Z' \) there is a scalar \( h_s \) with mass in the keV range.

To prevent possible reservations about these light states straightaway, let us discuss their coupling to the SM. The only way in which \( h_s \) couples to fermions other than neutrinos is via its mixing with the SM Higgs. Operators involving \( h_s \) linearly, thus, can be written as \( \mathcal{O}_{h_s} = c_{h_s} H S \times \mathcal{O}_{H S} \). Hence, couplings to fermions are suppressed by their Yukawa couplings and there are no new flavor changing effects. We adopt the bounds on this scenario of \([43]\). Besides BBN, which we discuss below, the strongest constraints arise from the burst duration of SN1987A and requires \( c_{H S} \approx 10^{-12} \) in the relevant parameter. Parametrically, \( (c_{H S} S_\nu S_\nu)^2 \approx \xi^2 \approx \varepsilon_m^2 \times 10^{-10} \), implying that we easily avoid this constraint for \( \varepsilon_m \lesssim 0.1 \).

The dominant coupling of \( Z' \) to SM fermions other than neutrinos is by mixing to the Z. Given Eq. (17), \( Z' \) couples to the SM neutral current with strength

\[
2 g_X (v_\phi/v_h)^2 = 2 g_X \xi^2 s_{\gamma}^2. \tag{26}
\]

For momentum transfer below \( m_Z \), this gives rise to new four-fermi (and NSI) operators of effective strengths

\[
\left( G_{\text{eff}}^{(2)f \nu S} \right)^2 / G_F = -2 \sqrt{2} \varepsilon_m s_\gamma^2, \quad \text{and} \quad \left( G_{\text{eff}}^{(4)f \nu S} \right)^2 / G_F = 4 \sqrt{2} \xi^2 s_\gamma^4 \approx \varepsilon_m^4 s_\gamma^4 \times 2 \times 10^{-9}. \tag{30b}
\]

Such feeble effects are currently not constrained by experiment.
We note that vector mixing can be modified by gauge-kinetic mixing of the U(1) field strengths $\mathcal{L}_\chi = -(s_\chi/2)B^{\mu\nu}X_{\mu \nu}$ [44, 45]. This shifts the $Z'$ coupling to the SM neutral current by a negligible amount proportional to $\chi \mathcal{O}(m_{Z'}^2/m_{T}^2)$ [46, 47] (given $m_{Z'} \ll m_{Z}, \chi \ll 1$). More important is the introduction of a coupling of $Z'$ to the electromagnetic current scaling as $c_X e \chi X$. Experimental constraints on this are collected in [48, 49] and our model could, in principle, saturate these limits. Therefore, we stress that $\chi \neq 0$ would neither affect our solution to the Hubble tension, nor the $H$ and $Z$ decay rates in Eqs. (33) and (34) below (to leading order), which are fixed by Goldstone boson equivalence.

We thus shift our attention to effects directly involving neutrinos. For $T \lesssim v_\phi$, neutrino mixing as in (19) is active. As required by direct-search bounds [50–52] and PMNS unitarity [53, 54] we are assuming [55, 56]

$$
e^m \leq 0.050, \quad (e^m) \leq 0.021, \quad (\tau^m) \leq 0.075, \quad \text{(31)}$$

for the mixing with $e, \mu, \tau$ flavors. The couplings of neutrinos to $Z'$ at low $T$ then are given by $g_X e^m Z'_\mu \chi$, with a strength set by (4). This gives rise to the four-fermion operators (2,30), but also to the possibility of $Z'$ emission in processes involving neutrinos. We stress that (4), together with $g_X \lesssim 1$, gives rise to a lower bound $e^m \gtrsim 5 \times 10^{-4}$, only two orders of magnitude below the limits (31). This fuels the intuition that this model is testable.

Constraints on neutrinos directly interacting with light mediators are collected in [21, 57–62]. The strongest laboratory constraints arise from meson [57, 63, 64] and nuclear double-beta decays [65–68]. However, even the most stringent bounds for the least favorable choice of flavor structure do not exclude couplings $g_{eff} \lesssim 10^{-5}$ for light $m_{Z'}$, comfortably allowing (4). While most of the laboratory constraints are interpreted in terms of light scalar (majaron) emission, the present study makes it worthwhile to revisit experimental exclusions in this region also for light vectors. The most important constraint is SN1987A neutrino propagation through the cosmic neutrino background (CvB) [69]. The exact bound depends on the neutrino masses and rank of $y$, but even under the most pessimistic assumptions $g_{eff} \lesssim 5 \times 10^{-4}$ cannot be excluded for $m_{Z'} \lesssim 60$ eV.

The $\nu\bar{\nu} \leftrightarrow \bar{\nu}\nu$ scattering cross section via $Z'$ exchange is approximately given by

$$
\sigma(\nu\bar{\nu}) = \frac{g_X^2 e^m}{12 \pi} \left( \frac{m_{Z'}^2 - s^2}{m_{Z'}^2} \right)^2 + m_{Z'}^2 \Gamma_{Z'}^2, \quad \text{(32)}
$$

To obtain the interaction rate in Fig. 1 we include the $t$-channel and use Maxwell-Boltzmann thermal average [70], while noting that a more refined analysis should employ Fermi-Dirac statistics [71, 72]. For $m_{Z'} > 2m_N$, $Z'$ decays to $N_1 N_1, N_2 N_2, \bar{\nu}N_2 (\nu N_2)$ and $\nu \bar{\nu}$, while for $m_{Z'} \lesssim m_N$ only the last channel is accessible. The respective total widths are $\Gamma_{Z'}/m_{Z'} \approx 10^{-7}$ or $10^{-12}$, see Fig. 1, corresponding to $Z'$ lifetimes from micro- to tens of picoseconds. For temperatures $T \gg v_\phi$ a thermal QFT investigation becomes necessary. We show thermally averaged rates (dashed) as obtained by dimensional analysis for illustration. For temperatures $T \ll m_{Z'}$ we reproduce the scaling of the effective operator (1) (also dashed). Before recombination, $\Gamma_{eff}(T)$ differs from the effective theory. While this should not change conclusions based on the (non-)free-streaming of neutrinos [24–29], see also [73], it certainly motivates dedicated cosmological analyses to tell if our specific temperature dependence could be discriminated from the effective model.

Finally we discuss the coupling of neutrinos to the light scalar $h_S$. Note that the matrix (19) is diagonalized exactly in $s_m$, reflecting massless active neutrinos and prevailing lepton number conservation at this stage. This prevents a quadratic coupling of neutrinos to $h_S$. Hence, SM neutrinos couple to $h_S$ only in association with hidden neutrinos, or suppressed by their tiny mass (e.g. $[M_{\nu}]_{11} \sim m_{\nu}$ produces such a coupling). In both cases effects are observably small, also because of the vastly suppressed coupling of $h_S$ to matter targets.

Also modification of $Z$ decays to neutrinos is observably small. Even if hidden $N$ sizably mixes with $\nu$, the invisible $Z$ width is not affected for $m_N \ll m_2 [50]$. However, the vertex $Z N \nu$ leads to $N$ production from neutrino upscattering on matter targets. The relevant operator is suppressed by $e^m$ compared to $G_F$. While interesting per se, $N$ decays invisibly, leaving an unaccompanied recoil as only signature.

Any consistent model of strong neutrino self- interactions requires a modification of the SM scalar sector and these are amongst the most visible effects of this model. The necessary modifications allow for new exotic decays of the SM $Z$ and Higgs bosons to invisible final states. To leading order in $\xi$ the rates of the most prominent decays are

$$
\Gamma_{H \rightarrow h_S h_S} = \frac{v_h^2}{32 \pi m_H} \left[ \lambda_{HS} e^2 + \lambda_{34} s^2 - \frac{\mu_{s2s}}{v_h} \right]^2, \quad \text{(33)}
$$

$$
\Gamma_{H \rightarrow Z' h_S} = \Gamma_{H \rightarrow h_S h_S}, \quad \Gamma_{Z' \rightarrow h_S} = \frac{m_Z g_{34}^2 s^4}{192 \pi c_W^3}, \quad \text{(34)}
$$

and to leading order in $\xi$ and $\gamma$

$$
\Gamma_{H \rightarrow ZZ'} \approx \frac{g_{34}^2}{c_W} \left( \frac{m_H - m_{Z'}^2}{m_H m_{Z'}^2} \right)^{\frac{3}{2}} \frac{\xi^2 s^4}{16 \pi} \left( 1 + \frac{\lambda_{HS}}{4 \lambda_H} \right)^2, \quad \text{(35)}
$$

Using $\Gamma_{H \rightarrow inv.} \leq 1.3$ MeV [74–76] and $\Gamma_{new Z' inv.} \leq 2.0$ MeV [77] we obtain constraints on the parameter space as shown in Fig. 2. $\Gamma_{Z' \rightarrow h_S}$ requires $\xi \lesssim 0.4$, while $\Gamma_{H \rightarrow inv.}$, in the absence of fine tuning, demands $\lambda_{HS}, \lambda_{34}, (\mu_{s2s}/v_h) \lesssim O(10^{-2})$. In the light of this, $\Gamma_{H \rightarrow ZZ'}$ is merely a rare Higgs decay with BR($H \rightarrow ZZ'$) $\approx 10^{-8} e^m$. A model similar to ours but with $S$ removed is phenomenologically excluded by $\Gamma_{Z' \rightarrow h_S h_S}$, which would have a rate as in (34) with $\gamma \rightarrow \pi/4$.

The charged scalars $H^\pm$ couple directly to charged leptons and hidden neutrinos via (8), with strengths set by $y$. Important constraints on $y$ arise from $\ell_1 \rightarrow \ell_2 \gamma$ and the measured lepton magnetic moments, both mediated by a loop of $H^\pm$ and $N$. Exact constraints are given in [55], while here it suffices to note that certainly $y \lesssim O(1)$ for all flavors, as we will find much tighter constraints below.
parameters have been chosen as $\gamma = 0.05$, $g_X = 2 \times 10^{-3}$, $\lambda_{H\nu} = 0.001$, $\lambda_3 = 0.002$, $\lambda_4 = 0.003$, $\lambda_\Phi = 0.3$, $\lambda_\Sigma = 0.4$, and $\lambda_{S\Sigma} = 0.5$.

The coupling of $H^\pm$ to quarks is suppressed by $s_\Sigma$ such that standard LHC searches [78, 79] do not apply. At LEP, $H^\pm$ could have been pair-produced via $s$-channel $\gamma/Z$ or $t$-channel hidden neutrinos, or singly-produced in association with charged and neutral leptons. $H^\pm$ dominantly decays to final states $N_\alpha \ell_\beta$ with BRs set by $y$. $N$ further decays to three neutrinos via $Z'$. The final state for $H^\pm$ hence is $\ell_\beta + \text{MET}$. We use LEP limits on $H^\pm$ pair-production [80] as well as a reinterpreted LEP selectron search [81–84] to obtain a lower bound $m_{H^\pm} > 100 \text{ GeV}$ [106].

Regarding electroweak precision tests, there are no new tree-level contributions to the $\rho \equiv O(T)$ parameter as we only introduce EW doublets and singlets. We follow [85] to estimate one-loop corrections. $T$ is always enhanced compared to the SM one-loop contribution, and stays in the allowed interval $T = 0.09 \pm 0.13$ [86] for $|m_{H^\pm} - m_{\Phi} | \lesssim 120 \text{ GeV}$. For $\lambda_4 > 0$ the new scalars $\Phi$ and $A$ are heavier than $H^\pm$. We assume them to be heavier than $m_{H^\pm}$ to avoid a small parameter window with mixed heavy-light decays.

**BBN.** —With $Z'$, $N$, and $h_S$ there are three new light species which could potentially distort BBN. Ultimately, a thermal QFT analysis seems worthwhile to fully explore the early universe cosmology of this model for $T \gg v_\nu$. The full set of coupled Boltzmann equations then should be solved to track abundances precisely, but this is beyond the scope of this letter. Nonetheless, a simple order of magnitude estimate suffices to clarify that there is a parameter region in which BBN can proceed as usual.

The tight bound on $\Delta N_{\text{eff}}$ during BBN [32] does not allow any of the new particles to be in thermal equilibrium with the SM. While $m^2_\nu$ is fixed by (5), the mass of $N$ is limited by (22), and $m_{h_S} \approx \xi_{v_N} \sqrt{2} \lambda_S \approx 7 \times 10^2 \sqrt{\lambda_S} \text{ MeV}$. Hence, given the allowed parameter space, neither of these states can simply be pushed beyond the MeV scale in order to avoid BBN constraints. Instead, we discuss the possibility that all of the new states are sufficiently weakly coupled to the SM at the relevant temperatures such that a thermal abundance is not retained.

While all of the new fields thermalize with the SM at EW temperatures, this connection is lost once the heavy scalars freeze out and decay. The initial abundance of new states subsequently is depleted by reheating in the SM, for example, at the QCD phase transition. We thus focus on the temperature region around BBN. The coupling of $Z'$ to the SM, as well as active hidden neutrino mixing $\varepsilon_m$, is only effective after the $U(1)_X$ breaking phase transition. This warrants that $Z'$ and $N$ do not thermalize with the SM between EW and BBN, realizing a generic mechanism [87] to reconcile short baseline neutrino oscillation anomalies with cosmology. Nonetheless, $Z'$ exchange would thermalize an abundance of $N$'s, which thus has to be absent. The leading process thermalizing $N$’s with the SM is $e^+e^- \leftrightarrow N\bar{N}$ via $t$-channel $H^\pm (\Phi, A)$ exchange, which scales as $\Gamma \sim (y/m_{H^\pm(\Phi)})^4 T^5$. Requiring this to be absent after QCD (EW) epoch requires $y \lesssim 6 \times 10^{-3.5} (m_{H^\pm(\Phi)}/100 \text{ GeV})$. Together with above bounds on $\varepsilon_m$ and $\gamma \lesssim 0.2$ this implies an upper limit $M \lesssim 300 \text{ eV}$ (QCD), or $M \lesssim 3 \text{ eV}$ (EW). For consistency of our analysis we have to require $m_\nu \ll y v_\nu \ll M$, implying a lower bound $y \gg 2 \times 10^{-5} (m_{H_S}/0.05 \text{ eV})$. Noteworthy, this allows $M_N$ right in the correct ballpark to resolve short baseline neutrino oscillation anomalies; not only in the well-known way with eV-scale states, see e.g. [88, 89], but also providing a definite model realization to the idea of decaying sterile neutrino solutions [90, 91].

The most relevant processes for thermalization of $h_S$ are $e^+e^- \leftrightarrow h_S h_S$, $e^-\gamma \leftrightarrow e^- h_S$, and $\nu\bar{\nu} \leftrightarrow h_S h_S$. None of them ever reaches thermal equilibrium due to the highly suppressed couplings of $h_S$.

Finally, we note that despite bearing some danger for successful BBN, the new states $N$ and $h_S$ can also be a virtue: In order to explain the Hubble tension with self-interacting neutrinos, $N_{\text{eff}}$ must be enhanced to $\Delta N_{\text{eff}} \approx 1$ during recombination [28, 29], requiring some energy injection in the dark sector after BBN [107]. In particular, having $m_{h_S} \sim O(10 \text{ keV})$ implies that an $h_S$ abundance could be present in a non-thermal state during BBN, and subsequently decay to reheat the neutrino background. $h_S$ decays to $\bar{N}\nu$ and $\bar{N}N_2$, but in absence of $L$-violation not to $\nu N_2$, $\bar{N}(1/2)N_1(2)$ or $\bar{\nu}\nu$ (or all processes barred), with a total width proportional to $\Gamma_{h_S} \propto y^2 s_\nu^2$, $\Gamma_{h_S}$, therefore, is extremely dependent on the exact parameters ranging somewhere from milli- to picoseconds. Also two-body decays $N \rightarrow Z'\nu$...
from a supposed non-thermal population of hidden neutrinos could contribute to $\Delta N_{\text{eff}}$ during CMB, provided $M_N > m_{Z'} + m_{\nu}$. For $M_N < m_{Z'} + \nu$, on the other hand, only three body decays $N \rightarrow (2\nu)\nu$ are possible with lifetime scaling as $\tau_N \sim (8\pi)^{-3}M_N^3 \epsilon_{\text{eff}}^2 \sim e^{-2}$. Depending on the exact parameters, a population of $N$, thus, could but doesn’t have to decay before recombination.

**Discussion.**—In summary, we have presented a consistent (renormalizable and phenomenologically viable) model that leads to vector-mediated neutrino self-interactions. In a narrow region of parameter space these interactions have the right strength to resolve the tensions between local and global determinations of $H_0$ and $\sigma_8$ [108]. To consistently implement such interactions in the SM, we had to introduce a second Higgs doublet and a hidden Dirac neutrino, both charged under a new gauge symmetry $U(1)_X$. Phenomenological consistency (invis-

The scalar spectrum then consists of several new states, all with very lepton-specific couplings: $h_{\nu}$ with mass of $O(10 \text{ keV})$, as well as $\Phi$, pseudo-scalar $A$ and the charged scalars $H^\pm$ all with masses of $O(100 \text{ GeV})$. The new, naturally neutrophilic fore carrier has a mass of $m_{Z'} \sim O(10 \text{ eV})$ and the new hidden neutrinos masses in the range $M_N \sim 1 \div 300 \text{ eV}$. A preferred region of parameter space has charged Higgses at a few 100 GeV, sizable $\text{BR}(\text{Higgs} \rightarrow \nu)$ as well as eV-scale hidden neutrinos. Other, perhaps testable signatures would then be non-standard neutrino matter interactions of strength $\epsilon_{(2\nu)}^{(2\nu,\nu)} \sim O(10^{-4})G_F$, cf. Eq. (30a), as well as active-hidden neutrino mixing with an angle $\epsilon_{\nu} > 5 \times 10^{-4}$.

That our model works without specifying the mechanism of neutrino mass generation may feel like a drawback to some. However, we think it is a virtue, as it renders this scenario compatible with all standard neutrino mass generation mechanisms.

The least appealing feature of our model, perhaps, is the introduction of several new scales $(v_\nu, v_8, \mu, M)$, and some hierarchies among them. We have nothing to say here about this or any other hierarchy problem but simply accepted this fact for the reason that we are convinced that this is the simplest renormalizable model in which active neutrinos pick up gauged self-interactions. Stabilizing these hierarchies against radiative corrections might require smaller scalar quartic cross-couplings than the direct constraints discussed above. Suchlike would not contradict any of our findings.

Finally, our analysis also shows that “model independent” considerations, which previously seemingly ruled out this model, are actually not always valid in concrete models. On the contrary, it is only in complete and consistent models that early universe cosmology like the Hubble tension can, and in fact must be, directly related to physics testable in laboratories.

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**Note added.**—During the completion of this work, Ref. [96] appeared on the arXiv which discusses a different UV complete model for self-interacting neutrinos. Also Ref. [97] appeared very recently which discusses bounds on self-interacting neutrinos from an EFT perspective. The corresponding bounds, where applicable, do not constrain our region of parameters.

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A model very similar to ours, albeit in a very different region of parameter space, has been investigated in [55, 56, 98]. However, baryons are neutral under our U(1)X, implying the absence of large non-standard neutrino-matter interactions.

We use abbreviations $\sin \theta_i \equiv s_i$, $\cos \theta_i \equiv c_i$, and $\tan \theta_i \equiv t_i$ for all angles $\theta_i$ in this work.

$s_X$ can be modified by gauge-kinetic mixing, we comment on this below Eq. (30).

We use the tadpole conditions to trade $\mu_H, \phi, s$ for other parameters. $m_{\mu_X}^2, m_{\phi}^2 > 0$ imply constraints on the allowed parameter space.

Sometimes a bound of $m_{\phi} > 275$ GeV is quoted based on [99]. However, this crucially depends on assumptions of $H^\pm$ BRs see [100, Sec. 4.7] for details.

As pointed out to us by S. Vogl, the thermalization of $N, h, S, Z'$ with $n$ flavors of active neutrinos after BBN produces entropy, which is released back to the neutrino background after the new states decay, before recombination. This would heat the neutrinos to $T_n' \approx [1 + 30/\tau_n]^{1/3}T_\nu$, cf. [101], giving rise to $\Delta N_{\text{eff}} \approx 1.03, 0.90, 0.74$ for $n = 3, 2, 1$.

We remark that the model is phenomenologically viable also for a different parameter region not discussed in this letter: For $m_{Z'} \ll eV$ and gauge coupling strength $g_X \ll 10^{-7}$. active neutrinos would only recouple after recombination with crucial impacts on the CRb.