The EDGES signal: An imprint from the mirror world?

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Recent results from the Experiment to Detect the Global Epoch of Reionization Signature (EDGES) show an anomalous spectral feature at redshifts $z \sim 15 - 20$ in its 21-cm absorption signal. This deviation from cosmological predictions can be understood as a consequence of physics that either lower the hydrogen spin temperature or increases the radiation temperature through the injection of soft photons in the bath. In the latter case, standard model neutrino decays $\nu_i \rightarrow \nu_j \gamma$ induced by effective magnetic and electric transition moments ($\mu_{\text{eff}}$) are precluded by the tight astrophysical constraints on $\mu_{\text{eff}}$. We show that if mirror neutrinos are present in the bath at early times, an analogous mechanism in the mirror sector can lead to a population of mirror photons that are then “processed” into visible photons through resonant conversion, thus accounting for the EDGES signal. We point out that the mechanism can work for mirror neutrinos which are either heavier than or degenerate with the standard model (SM) neutrinos, a scenario naturally realized in mirror twin Higgs models.

I. INTRODUCTION

After recombination the universe was filled with radiation, dark matter (DM) particles and primordial gas (mainly hydrogen). This cosmic stage, known as the “dark ages”, lasted until the formation of the first structures, an event that started when Compton scattering processes could not maintain the gas and the radiation in equilibrium. The gas, being cooled faster than the radiation field, got gravitationally trapped in DM haloes and eventually ended up collapsing and fragmenting, giving rise to the appearance of stars, quasars and galaxies. At lower redshifts the Lyman-alpha photons emitted by these first structures led to a re-ionization period, known as the re-ionization era.

The only known observable with which the dark ages can be observationally accessed is the 21-cm line of the ground-state hyperfine transition of atomic hydrogen [1]. This probe provides as well a way to test the re-ionization epoch, thus allowing the study of the cosmic time when astrophysical objects became the dominant source of the intergalactic medium. The cosmic microwave background (CMB) photons are resonantly absorbed by the hydrogen atoms, thus producing a change in the radiation field, got gravitationally trapped in DM haloes and eventually ended up collapsing and fragmenting, giving rise to the appearance of stars, quasars and galaxies. At lower redshifts, the Lyman-alpha photons emitted by these first structures led to a re-ionization period, known as the re-ionization era.

The result of the recoupling of $T_{\text{gas}}$ and $T_s$ due to Lyman-alpha photons from early stars. The observed absorption profile is centered at around $z \sim 17$ and covers redshifts in the range 15-20 [6]. To a large extent, the profile is consistent with cosmological predictions, but the observed amplitude indicates more absorption than expected. The, arguably, most simple explanation would be an earlier $T_{\text{CMB}} - T_{\text{gas}}$ decoupling (at $z \sim 250$ rather than $z \sim 150$), which would produce an earlier cooling of the gas. This, however, does not work since it requires the ionization fraction to be less than the expected fraction by about an order of magnitude, something strongly disfavored by Planck data [6].

An explanation of the observed spectral profile requires either decreasing $T_{\text{gas}}$ (gas cooling) or increasing $T_{\text{CMB}}$ (radiation heating). And indeed since the release of the EDGES result both alternatives have been studied in the literature. Ref. [7] considered DM-baryon scatterings determined by a velocity-dependent cross section resulting from a Coulomb-like interaction. After discarding the possibility of a light mediator due to fifth force constraints, ref. [8] showed that subdominant millicharged DM can explain the 21-cm spectral feature, despite in a constrained region in parameter space that must be endowed with an additional depletion mechanism to prevent overproduction. It has been pointed out that this constraints can be relaxed provided the millicharged DM is produced after recombination [9]. Ref. [10] discussed various mechanisms, among which those based on the emission of soft photons that can heat up the radiation temperature [11]. Using dipole DM as a benchmark model [12], it ruled out these kind of scenarios. Other mechanisms put forward include black hole remnants from Pop-III stars [13], interacting dark energy models [14], charge sequestration models [15] and more relevantly for our study dark-photon to photon resonant conversion [16]. This latter relies on a non-thermal population of dark photons, resulting from the decay of an unstable relic, which are then resonantly converted into photons at redshifts $z \sim 17$.

Neutrinos can couple to electromagnetic radiation through electric charge (milli-charged), electric/magnetic dipole (trans-
we consider mirror neutrino decays and resonant conversion \( \mu \) the photon flux assuming the redshifts relevant for EDGES, we discuss in more detail the conditions required for addressing the EDGES signal. In this paper, we start by checking whether these bounds one could moderately raise the radiation temperature by injecting photons though neutrino decays. After showing that the bounds on \( \mu_{\text{eff}} \) always lead to a suppressed photon flux, we then entertain the possibility that mirror neutrinos endowed with the same type of couplings can inject a sufficiently high photon flux so to enable addressing the EDGES anomalous spectral feature. We study in detail mirror neutrino decays to mirror (dark) photons, \( \nu_i' \rightarrow \nu_j' + \gamma' \), occurring at high red-shift and then getting resonantly converted into visible photons \( \gamma' \rightarrow \gamma \). For that aim we consider two scenarios inspired in mirror twin Higgs models defined by degenerate and non-degenerate SM and mirror neutrino masses with \( T' < T \) (where \( T' \) and \( T \) refer to the mirror and SM temperatures respectively), as required by cosmological constraints on additional dark radiation \( \Delta N_{\text{eff}} \) [19].

The rest of the paper is organized as follows. In sec. [II] we discuss generalities on the 21-cm absorption signal and settle the conditions required for addressing the EDGES signal. In sec. [III] we consider the case of SM neutrino decays during the redshifts relevant for EDGES, we discuss in more detail current bounds on neutrino transition moments and calculate the photon flux assuming \( \mu_{\text{eff}} \) is a free parameter. In sec. [IV] we consider mirror neutrino decays and resonant conversion of dark photons into visible ones. We then provide a theoretical motivation in sec. [V] based on mirror twin Higgs models, for the mirror neutrino scenarios we consider. In sec. [VI] we summarize and present our conclusions.

II. GENERALITIES

During the recombination era \( (z \sim 1100) \) electrons and protons recombined to form neutral hydrogen. As shown by the high degree of isotropy of the CMB, the universe was highly uniform at that time thus suggesting that few, if any, luminous objects could have formed. Adiabatic expansion thus led to a stage in which the universe consisted mainly of a neutral gas, CMB photons and DM particles, a cosmic stage known as the dark age. The universe evolved adiabatically and the radiation temperature, \( T_{\text{CMB}} \), decreased with redshift according to \( T_{\text{CMB}} = 2.7(1+z) \) K. The remaining small ionization fraction, \( X_e = n_e / n \), enabled the injection of energy from the CMB to the gas through Compton scattering processes, thus keeping both baryons and radiation at the same temperature until \( z \sim 150 \).

The virial temperature of a DM halo \( (T_{\text{vir}}) \) determines the binding energy of the material within the halo. Accordingly, only gas for which \( T_{\text{gas}} < T_{\text{vir}} \) can be trapped by the halo gravitational pull. For \( z \lesssim 150 \), Compton scattering effects became less effective and so the temperature of the gas decreased faster than the radiation temperature. The gas then was trapped by the DM halo, but the shocks induced by the gravitational collapse heated up the gas to \( T_{\text{vir}} \), thus driving the system to hydrostatic equilibrium. After departing from this state, the gas contracted within the halo and became gravitationally stable, at some point it fragmented and led to the formation of the first stars, quasars and galaxies. The high-energy radiation emitted from these first objects reionized the hydrogen in the intergalactic medium, leading to the re-ionization epoch.

The only known observable with which the dark age period can be studied is the redshifted hydrogen hyperfine transition spectral line. It enables as well detailed studies of the epoch of re-ionization such as structure formation and the formation of the first galaxies. The ground state of neutral hydrogen is split into two hyperfine states due to proton-electron spin-spin coupling: a singlet, corresponding to the anti-alignment of the two spins and a degenerate triplet state corresponding to the alignment of both spins. The energy splitting between these states is \( \Delta E = E_1 - E_0 \approx 5.9 \mu eV \), which corresponds to a \( \sim 21 \) cm photon wavelength and a rest-frame frequency \( \nu_{10} = 1420 \) MHz, redshifted as \( \nu(z) = 1420/(1+z) \) MHz. Some of the CMB photons propagating in the medium can be absorbed by hydrogen resulting in a singlet-triplet transition which modifies the brightness temperature of the CMB according to [20]

\[
T_b(z) = T_{\text{CMB}}(z)e^{-\tau(z)} + \left(1 - e^{-\tau(z)}\right)T_i(z).
\]

Here \( T_{\text{CMB}}(z) \) is the brightness temperature of the CMB without absorption, \( T_i(z) \) is the spin temperature which characterizes the relative population of the triplet to the singlet states. The optical depth reads

\[
\tau(z) = \frac{3e^2h_pA_{10}n_{\text{HI}}(z)}{32\pi\nu_{10}^2k_BT_i(z)H(z)},
\]

with \( A_{10} \approx 2.9 \times 10^{-15} \) s\(^{-1} \) the spontaneous decay rate for the excited to the ground hyperfine states, \( n_{\text{HI}} \) the density of neutral hydrogen, \( c \) the speed of light, \( h_p \) the Planck constant, \( k_B \) the Boltzmann constant and \( H(z) \) the Hubble expansion rate. Since \( \tau \ll 1 \), the change in the brightness temperature seen today \( T_{21}(z) = (T_b(z) - T_{\text{CMB}}(z))/(1+z) \) can be recast as follows

\[
T_{21} \simeq \mathcal{F}\left(\frac{0.15}{\Omega_mh^2}\right)^{1/2}(1+z)^{1/2}\left(\frac{\Omega_bh^2}{0.02}\right)^{1/2}\left[1 - \frac{T_{\text{CMB}}(z)}{T_i(z)}\right].
\]

where \( \mathcal{F} = 2.3 \) mK \( x_{\text{HI}}(z) \) (\( x_{\text{HI}} \) is the neutral hydrogen fraction), \( \Omega_m \) and \( \Omega_b \) are respectively the matter and baryon energy densities in units of the critical density and \( h \) is the Hubble constant in units of 100 km/s/Mpc.

At the center of the absorption profile the redshift amounts to \( z \sim 17 \). The quantities entering (3) at such redshift take the following values: \( x_{\text{HI}}(17) \approx 1 \), \( T_i(17) = T_{\text{gas}} \approx 7 \) K [21] and
\( T_{\text{CMB}} = 2.7 \times 18 \text{K} \approx 49 \text{K} \). Thus, the expected value for the brightness temperature contrast is \( T_{21}(z = 17) = -0.2 \text{K} \). The value provided by EDGES for the same redshift is in contrast \( T_{21}(z = 17)_{\text{EDGES}} = -0.5^{+0.5}_{-0.3} \text{K} \) at 99\% CL, which corresponds to about a 3.8\( \sigma \) deviation from theoretical expectations. In general \( T_{\text{gas}} \leq T_{s} \leq T_{\text{CMB}} \), and so the lowest spin temperature corresponds to the case \( T_{s} = T_{\text{gas}} \) (full Lyman-alpha coupling) \(^{20}\), thus representing the case where \( T_{\text{gas}} / T_{s} \) is the largest. Under this assumption, the EDGES signal can be reconciled if the ratio \( T_{\text{CMB}} / T_{s} \) is enhanced by a factor 2. If one departs from this assumption and considers the more general case where \( T_{s} > T_{\text{gas}} \) this factor should increase accordingly \(^{22}\). For concreteness throughout our analysis we will assume \( T_{s} = T_{\text{gas}} \).

### III. EXTRA RADIATION FROM ELECTROMAGNETIC-INDUCED NEUTRINO DECAYS

The injection of soft photons in the early universe by the SM neutrino decays can proceed through magnetic (electric) transition moments \( \mu_{ij} \left( \epsilon_{ij} \right) \). In the following we consider the effective electromagnetic neutrino interactions

\[
L_{\text{eff}} = \frac{1}{2} v_{i} \sigma_{\mu\nu} \left( \mu_{ij} + \epsilon_{ij} \gamma^{5} \right) v_{j} F_{\mu\nu},
\]

where \( \sigma_{\mu\nu} = i \left[ \gamma_{\mu}, \gamma_{\nu} \right] / 2 \). \( F_{\mu\nu} \) is the electromagnetic field strength tensor and \( i,j \) label neutrino mass eigenstates. Note that here we have assumed neutrinos are Dirac particles, assuming otherwise will not change our conclusion. These couplings induce radiative neutrino decays \( v_{i} \rightarrow v_{j} + \gamma \) for which the decay width can be written as

\[
\Gamma_{v_{i} \rightarrow v_{j} + \gamma} = \frac{\mu_{\text{eff},ij}^{2}}{8\pi} \left( \frac{\Delta m_{ij}^{2}}{m_{i}} \right)^{3},
\]

with \( \mu_{\text{eff},ij} = \sqrt{\left| \mu_{ij}^{2} \right| + \left| \epsilon_{ij} \right|^{2}} \). \( \Delta m_{ij}^{2} = m_{i}^{2} - m_{j}^{2} \) and \( m_{i} \) the i-th neutrino mass eigenstate. (Where family index is not relevant, neutrino mass will be denoted simply as \( m_{\nu} \).) This effective coupling is subject to tight constraints from laboratory experiments and astrophysical considerations (see e.g. \(^{17}\)). For the former, the most severe bound is derived from the GEMMA experiment which relies on measurements of electron recoils induced by the neutrino-electron elastic scattering process \( \nu e \rightarrow \nu e \) \(^{23}\). The current 90\% CL limit neglecting atomic effects reads (from now on we will drop family indices, except in those cases where strictly necessary)

\[
\mu_{\text{eff}} \lesssim 3.2 \times 10^{-11} \mu_{B}.
\]

Astrophysical bounds are more stringent, in particular those derived from plasmon decay (\( \gamma \rightarrow v v \)) in globular cluster stars \(^{18}\). This process—enabled by medium effects—releases an amount of energy through the neutrinos that escape the stellar medium, resulting in a delay in helium ignition and thus in the following upper limit

\[
\mu_{\text{eff}} \lesssim 3.0 \times 10^{-12} \mu_{B}.
\]

Next, notice that photons injected much before recombination \( (z \sim 1100) \) will get fully absorbed by the plasma, while those injected below \( z \sim 15 \) cannot contribute to the spectral distortion observed by EDGES. Thus, if photons emitted in neutrino decays were to be responsible for the EDGES signal, they should be generated in the window \( 15 \lesssim z \lesssim 1100 \) with energy falling within the EDGES energy absorption interval \((0.28, 0.37) \mu \text{eV} \) (rest-frame frequency redshifted in the interval \( 15 < z < 20 \)). To determine which photons can contribute to the signal, one needs their energy at production properly redshifted, \( c'\text{-est-à-dire} E(z) = \Delta m_{ij}^{2}/2/m_{i}/(1+z) \), assuming decays at rest. The condition of this energy falling within the absorption energy range, fixes the minimum (and maximum) mass that the decaying neutrinos should have so to be “visible”. Using the best fit point values for the mass squared differences \(^{24}\), we find \( m_{\nu} \gtrsim 3 \text{eV} \) \( (h = 3 \text{ for normal order neutrino mass spectrum, } h = 1 \text{ for inverted order}) \) and \( m_{3} \gtrsim 0.1 \text{eV} \). Cosmological constraints on neutrino masses \( \Sigma m_{\nu} \lesssim 0.68 \text{eV} \) (95\% CL limit) \(^{19}\), thus imply that photons produced by the decay of \( \nu_{e} \) will fall outside the EDGES energy window and only decay of \( \nu_{2} \) matters. To determine the contribution of neutrino decays to the number of photons in the plasma, one should calculate the photon number density per-unit energy, which at present time reads \(^{25}\)

\[
\frac{dn_{\gamma}}{dE} = \frac{B}{E} \frac{n_{\nu}(t_{0})}{H(z)} e^{-\left(\Gamma_{\nu}/H(z)\right)},
\]

where \( E \) is the photon energy today, \( B \) the branching fraction for the radiative decay, and \( n_{\nu}(t_{0}) \) the would-be present number density of neutrinos if they did not decay. Here \( \left\langle \Gamma_{\nu} \right\rangle \) is the thermally averaged total decay width of neutrino \( \nu \)

\[
\left\langle \Gamma_{\nu} \right\rangle = \frac{1}{3} \frac{K_{1} \left( mv_{h}/T_{\text{CMB}} \right)}{K_{2} \left( mv_{h}/T_{\text{CMB}} \right)},
\]

with \( K_{0} \) the temperature-independent total decay width and \( K_{i}(x) \) the order \( i \)-th modified Bessel function of the second type. The expansion time \( t(z) \) is given in terms of the Hubble expansion rate \( H(z) \), namely

\[
t(z) = \int_{z}^{\infty} \frac{dz'}{1+z'} \frac{1}{H(z')},
\]

which for a flat universe with matter \( \Omega_{m} \) and a non-vanishing cosmological constant \( \Omega_{\Lambda} \), is given by \(^{26}\)

\[
H(z) = H_{0} \sqrt{\Omega_{\Lambda} + \Omega_{m}(1+z)^{3} + \Omega_{\nu}(1+z)^{4}}.
\]

For our calculation we

\(^{2}\) We have included the radiation contribution \( \Omega_{r} = 5.38 \times 10^{-5} \) \(^{28}\) which is negligible for \( z \lesssim 1100 \) but will be important for our considerations in the next section.
have taken $H_0 = 67.8$ km/s/Mpc, $\Omega_\Lambda = 0.69$ and $\Omega_m = 0.31$ \cite{20}. Bearing in mind that the would-be number density of SM neutrinos per generation today is

$$n_\nu(t_0) = \frac{3}{2} \frac{\zeta(3)}{\pi^2} \left(\frac{4}{11}\right) T_0^3,$$  \hspace{1cm} (11)$$

where $T_0 = 2.725$ K is temperature of the CMB photons today, the contribution to the photon number density from $\nu_2$ decays can then be calculated. Assuming full Lyman-alpha coupling ($T_{\text{gas}} = T$) \cite{6}, this contribution should amount to that of the CMB so to account for the EDGES signal (see sec. \[11\]. Fig. \[1\] shows the result for the photon number density due to $\nu_2$ radiative decays in comparison to the Rayleigh-Jeans tail of the CMB black body spectrum. This result has been derived assuming a normal order neutrino mass spectrum (inverse mass ordering gives similar results), using $m_2 = 0.12$ eV which minimizes the required $\mu_{\text{eff}} = 7.8 \times 10^{-6} \mu_B$, the latter a value far larger than current limits. Thus, an explanation of the anomalous spectral distortion observed by EDGES, based on electromagnetic-induced radiative neutrino decays requires effective electromagnetic couplings already ruled out by data.

It is worth pointing out that even if one could afford a sufficiently large $\mu_{\text{eff}}$, there are extra effects one should deal with. First of all since $\Gamma_{\nu_1} > \Gamma_{\nu_2}$, $\nu_3$ decays will yield a more abundant photon flux in the energy range $\sim [2, 125] \mu$eV and will contribute sizeably to the CMB at redshifts above the EDGES window. That effect, however, could be kept under control by assuming that the effective transition magnetic moments of $\nu_3$ are suppressed. As can be seen in fig. \[1\] one finds the same effect for $\nu_2$. And of course in this case the solution used for $\nu_3$ will not work, implying that the scenario will be further constrained from measurements of distortions to the CMB at redshifts $z \gtrsim 50$.

IV. EXTRA RADIATION FROM A MIRROR SECTOR

The conclusion reached in the previous section might change if mirror neutrinos couple to radiation in the same way SM neutrinos do. Let us discuss this scenario in more detail. The electromagnetic couplings of the mirror neutrinos resemble those in $\text{(1)}$, with neutrinos and the electromagnetic field tensor traded for those of the mirror sector, which we will denote as $\nu'$ and $F_{\text{\nu'}}$. For the coupling we will use

$$\mu'_{\text{eff}} = \sqrt{|\mu_{\nu_{ij}}'|^2 + |e_{\nu_{ij}}'|^2},$$

with $\mu_{\nu_{ij}}'$ and $e_{\nu_{ij}}'$ the mirror neutrino magnetic and electric transition moments. In addition to these couplings one has as well a kinetic mixing term which couples the electromagnetic field tensors of the visible and mirror sectors, $(\epsilon/2)F^{\mu\nu}F'_{\mu\nu}$. The simultaneous presence of $\mu'_{\text{eff}}$ and $\epsilon$ induces processes of the type $\gamma \rightarrow \nu'\nu'$, which as in the SM case leads to stellar cooling and thus to the upper limit

$$\epsilon \mu'_{\text{eff}} \lesssim 3.0 \times 10^{-12} \mu_B.$$  \hspace{1cm} (12)$$

The mirror sector is subject as well to cosmological constraints which require the SM temperature to be larger than the mirror sector temperature. This can be understood from the contribution of mirror neutrinos to the effective “neutrino” degrees of freedom

$$\Delta N_{\text{eff}} = \frac{4}{7} \left(\frac{11}{4}\right)^{4/3} g_*' ,$$  \hspace{1cm} (13)$$

where $g_*$ refers to the effective relativistic degrees of freedom which is given by

$$g_*' = \sum_{i=\text{boson}} g_i' \left(\frac{T_i}{T_\gamma}\right)^4 + \frac{7}{8} \sum_{i=\text{fermion}} g_i' \left(\frac{T_i}{T_\gamma}\right)^4 .$$  \hspace{1cm} (14)$$

Here $T_\gamma$ is the photon temperature while $T_{\nu'}$ is the temperature of the corresponding mirror sector relativistic degree of freedom. Assuming $T_{\nu'}^i = T^i$ (common temperature for all mirror
sector relativistic degrees of freedom) and that at the time of \(\nu'\) decay only the dark photon and the three mirror neutrino species are relativistic, \(\Delta N_{\text{eff}}\) becomes

\[
\Delta N_{\text{eff}} = \frac{29}{7} \left( \frac{11}{4} \right)^{4/3} \left( \frac{T'_{\nu}}{T_q} \right)^4. \tag{15}
\]

Thus, by using the 2σ limit \(\Delta N_{\text{eff}} < 0.65\) from Planck (the bound from Big Bang nucleosynthesis (BBN) is comparable)\(^\text{[19]}\) an upper bound on the temperature of the mirror sector can be derived, \(T' < 0.45 T_q\). This bound could be relaxed if the heaviest and next-to-heaviest mirror neutrinos decay before BBN \((z \sim 4 \times 10^8)\) and the lightest one is stable, in that case the mirror sector can be slightly hotter, \(T' < 0.53 T_q\).

Another bound one has to consider has to do with mirror neutrino masses \(m_{\nu'} = \{m'_i\}\), which can be constrained by combining \(^\text{[15]}\) with cosmological limits on \(\sum m_i\) (see sec. \(\text{III}\)). Taking \(m_{\nu'} = r m_{\nu}\), with \(r\) a common rescaling that determines how heavy the mirror neutrinos can be, and using \(\sum m_i + \sum m'_i (n'_i/n_i) < 0.68\) eV \(^\text{[19]}\) we find the following upper limit

\[
r \lesssim 55 \left( \frac{0.65}{\Delta N_{\text{eff}}} \right)^{3/4} \left[ \left( \frac{0.05eV}{\sum m_i} \right) - \frac{5}{68} \right]. \tag{16}
\]

Here we normalize the sum of SM neutrino masses to the value of the atmospheric mass scale determined from neutrino oscillation data \(^\text{[24]}\).

### A. Dark photon resonant conversion

Mirror neutrino decays can directly generate a photon flux through kinetic mixing, \(\nu'_i \to \nu'_j + \gamma\). These decays however will be controlled by \(\varepsilon \mu_{\nu'}\), and so given the bound in \(^\text{(12)}\) the photon flux will be rather suppressed (pretty much resembling what we found in the SM neutrino case). On the other hand, mirror neutrinos decay to dark photons \((\nu'_i \to \nu'_j + \gamma')\) can yield a population which is not necessarily small. The key point is that these decays are solely determined by \(\mu_{\nu'}\), which can be large if \(\varepsilon \ll 1\) while satisfying \(^\text{(12)}\). These decays can take place way above \(z \sim 1100\), as far as they occur after the mirror and SM sectors have thermally decoupled, to avoid \(\gamma' - \gamma\) thermalization.

Once in the bath, as we will see shortly, depending on the mass of dark photons, they can be efficiently "processed" into visible photons through resonance conversion, even with small \(\varepsilon\) enforced by \(^\text{(12)}\) together with the CMB constraints \(^\text{[27]}\). In contrast to mirror neutrino decays, the conversion process should occur in the window \(15 \lesssim z \lesssim 1100\) if this mechanism is to explain the anomalous spectral profile reported by EDGES for the reasons elaborated after eq. \(^\text{(7)}\).

In the heat bath visible photons acquire an effective mass through the scattering with free electrons and neutral atoms. Neglecting the latter it can be written according to \(^\text{[27]}\)

\[
m_{\gamma} \simeq 1.75 \times 10^{-14} (1 + z)^{3/2} X_e^{1/2} \text{eV}, \tag{17}
\]

where the free electron fraction \(X_e\) can be well approximated for \(z \gtrsim 70\) by the expression \(^\text{[27]}\)

\[
\log X_e \simeq \frac{3.15}{1 + \varepsilon^2} \left( \frac{z}{z - 907}{160} \right). \tag{18}
\]

Resonant \(\gamma' - \gamma\) conversion, which resembles the MSW effect for solar neutrinos \(^\text{[28,30]}\), happens when the dark photon mass amounts that of the visible photon mass, \(m_{\gamma'} \simeq m_{\gamma}\). In that case the \(\gamma' - \gamma\) conversion probability can be taken as \(^\text{[27]}\)

\[
P_{\gamma' - \gamma} = P_{\gamma' - \gamma'} \simeq 1 - e^{-2\pi r k \sin^2 \varepsilon}, \tag{19}
\]

which holds for \(\varepsilon \ll 1\). The second term corresponds to the level crossing probability with \(k = m_{\gamma'}/(2E)/(1 + z)\) and \(r = \left| d \log m_{\gamma'}/dt \right|^{-1} \left| \frac{d \log m_{\gamma'}/dt}{1 + z}/H(z)\right|_{z = z_{\text{res}}}\).

The setup of eqs. \(^\text{(17)}\) and \(^\text{(18)}\) as well as the definitions for the parameters \(k\) and \(r\) allow the determination of \(P_{\gamma' - \gamma}\) and therefore of the corresponding photon spectrum

\[
\frac{dn_{\gamma}'}{dE} = P_{\gamma' - \gamma} \times \frac{B}{E} \frac{n_{\nu'} (t_0) (\Gamma'_{\gamma'})}{H(z)} e^{-\Gamma'_{\gamma'} t(z)} \Theta(z - z_{\text{res}}), \tag{20}
\]

where the Heaviside function assures that dark photons produced below \(z_{\text{res}}\) (the redshift for which the resonance occurs) will not be converted into visible photons. To show that this mechanism can account for the EDGES signal, we fix \(\mu_{\nu'} = 3 \times 10^{-5} \mu_B\), \(\varepsilon = 10^{-3}\) and \(T'/T_q = 0.4\). For illustration, we have chosen the redshifts for which resonance conversion occurs to be \(z_{\text{res}} = 1200\) and 1050. These fix \(m_{\gamma'} \simeq m_{\gamma}\) to be \(4.4 \times 10^{-10}\) eV and \(2.1 \times 10^{-10}\) eV respectively. The \(\varepsilon\) has been chosen such that it is consistent with the bounds from distortions of the CMB spectrum \(\varepsilon \lesssim 10^{-6}\). We then calculate the photon number density generated through \(\gamma' - \gamma\) conversion as a function of the heaviest mirror neutrino mass \(m_{\gamma'}\) assuming normal order. We include photons from mirror neutrino decays as well as from the dark background radiation \((n_{\nu'} = n_{\gamma'} + n_{\gamma'}^{\text{CMB}})\), and compare with the CMB photon number density (\(n_{\gamma'}^{\text{CMB}}\)). The result is shown in fig. \(^\text{(2)}\) where we have specified two scenarios for the mirror neutrino mass spectrum\(^\text{[14]}\) (a) complete degeneracy between SM and mirror neutrinos, (b) non-degeneracy, \(m_{\nu'} = r m_{\nu}\), with \(r\) subject to the bound in \(^\text{(16)}\). This result shows that one can address the EDGES spectral feature by means of this mechanism.

### V. REALIZATION IN THE MIRROR TWIN HIGGS MODEL

The mirror neutrinos we have considered in the previous section are naturally realized in twin Higgs models \(^\text{[32]}\). In a rather simple realization, one can understand

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\(^3\) Although we have chosen the normal order neutrino mass spectrum, inverse order does not change our results.
them as a class of models in which the scalar potential features a global $U(4)$ symmetry. The scalar field of the theory, $\mathcal{H}$, transforms as the fundamental representation of the global symmetry group and acquires a vacuum expectation value (vev) $v'$ that spontaneously breaks $U(4)$ to $U(3)$ leaving behind—in the absence of quantum corrections—seven Nambu-Goldstone bosons (NGBs). $\mathcal{H}$ is constructed out of two doublets belonging to the gauged direct product subgroup $SU(2) \times SU(2) \subset U(4)$, $H$ and $H'$ (one would identify $SU(2)$ with the SM $SU(2)_L$ and $H$ with the SM Higgs doublet). Since the global symmetry is explicitly broken by gauge $SU(2) \times SU(2) \subset U(4)$, one would expect the gauge quantum corrections to lift the mass of the NGBs. On the one hand, the presence of a $Z_2$ “twin” symmetry (that interchanges $H \leftrightarrow H'$) leads to self-energy one-loop corrections that are $U(4)$ invariant, and so do not contribute to $m_{\text{NGB}}$. On the other hand, gauge one-loop quantum corrections to the scalars ($H$ and $H'$) four-point functions, instead, explicitly break $U(4)$, thus implying $m_{\text{NGB}} \neq 0$ (the NGBs are actually pseudo-NGBs). These corrections, however, are logarithmically divergent and so allow $O(m_{\text{NGB}}) \sim g^2 v'/4\pi$ to be at the weak scale for cut-off scales up to $\sim 5$-10 TeV. Thus, identifying the SM Higgs among these degrees of freedom prevents the Higgs mass from acquiring large quantum corrections and therefore solves the little hierarchy problem, with new particles which are singlets under the SM.

In the limit of exact $Z_2$ symmetry both the SM Higgs vev, $v$, and $v'$ are equal. A mechanism that enables a mild hierarchy between $v$ and $v'$ relies on the introduction of a term in the scalar potential that softly breaks $Z_2$ and leads to $v < v'$. \textsuperscript{[32, 33]} This small $Z_2$ breaking is required to obtain the correct electroweak breaking scale and Higgs precision measurements further require $v'/v \gtrsim 3$. \textsuperscript{[34]} Complete models can be constructed with the aid of this mechanism, by extending the symmetry to all the interactions of the SM or by identifying the twin symmetry with parity, in which case two models are possible: mirror twin Higgs models \textsuperscript{[32]} or left-right symmetric twin Higgs models \textsuperscript{[35]}. Is within the former—in which there is a mirror copy of the SM with the same (mirror) particle content and interactions—that the scenarios we have pointed out in the previous section emerge and what we will focus on next. (Here $Z_2$ is the symmetry which interchanges between particles and mirror particles.)

In the mirror twin Higgs models, further call for $Z_2$ breaking arise from the constraints on extra radiation generated by the mirror sector \textsuperscript{[34]}. At $T > \text{GeV}$, both the SM and mirror sectors are thermally coupled through Higgs-exchanged processes. As the temperature decreases, mirror particles inject dark radiation in the bath (in the form of light degrees of freedom) through decay and annihilation processes. Their contribution to dark radiation is determined by the decoupling temperature $T_d$, below which SM and mirror sector interactions are slower than the Hubble expansion rate and both sectors decouple.

In the case of $Z_2$-symmetric Yukawa couplings, fermion masses in both sectors differ only by the ratio $v/v'$. And

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig2}
\caption{Extra radiation from mirror neutrino decay to dark photons and then processed to visible radiation through $\gamma' - \gamma$ conversion as a function of $R = (m_{\nu'}/0.05\text{eV})$ (assuming a normal order mass spectrum). The temperature of the mirror sector has been fixed to $T' = 0.4 T_d$ and $n_{\gamma} = n_{\gamma}^{\text{CMB}} + n_{\gamma}^{\text{decay}}$ ($n_{\gamma}^{\text{CMB}}$ refers to the amount of dark CMB radiation converted into visible radiation). We have chosen two benchmark values of the redshift for which resonant conversion occurs at $z_{\text{res}} = 1200$ and 1050 which correspond to $m_{\nu'}/m_{\nu} \approx 4.4 \times 10^{-10} \text{eV}$ and $2.1 \times 10^{-10} \text{eV}$, respectively. The vertical line which separates the two shaded regions corresponds to $\sum_i m_i + \sum_i m_i'(n_{\nu'}/n_{\nu}) = 0.68$ eV and in the left (right) region, the SM and mirror neutrinos can (cannot) be degenerate.
}
\end{figure}

\begin{itemize}
\item \textsuperscript{4} In fact, the scalar potential has an enhanced global symmetry $O(8)$ and one can consider the breaking as $O(8) \rightarrow O(7)$ which can contain the custodial symmetry of the SM.
\item \textsuperscript{5} This is a consequence of $g = g'$ (with $g, g'$ the $SU(2) \times SU(2)'$ gauge couplings respectively), implied by the discrete twin symmetry. Ultimately due to this symmetry the NGBs are insensitive to quadratic divergences.
\item \textsuperscript{6} Another way out is to remove the troublesome light mirror particles i.e. by having an imperfect copy of the SM in the mirror sector \textsuperscript{[35]}.
\end{itemize}
generated from dimension five effective operators, with which is consistent with the bound on \( \Delta N_{\text{eff}} \) demands this ratio to be rather large \( v/v' \gtrsim 40 \), implying a fine-tuning to get the correct electroweak scale \[34\]. Breaking \( \mathbb{Z}_2 \) in the Yukawa sector allows for \( y_F > y_f \) and so for heavier mirror charged leptons and quark masses which in particular imply higher mirror QCD phase transition temperature \( T'_{\text{QCD}} \) (up to \( \sim 3 \text{GeV} \)) \[33\]. The amount of dark radiation can then be reduced by having \( T_q \) below \( T'_{\text{QCD}} \) (so light mirror quarks will not contribute at decoupling) and above the SM one \( T_{\text{QCD}} \sim 0.2 \text{GeV} \).4 Assuming separate entropy conservation in the mirror and SM sectors below \( T_q \), their temperatures are related through \[38\]

\[
T' = \left[ \frac{g_*(T') g_*(T_d)}{g_*(T) g_*(T_d)} \right]^{1/3} T .
\]

For instance, assuming that at \( T_d \sim \text{GeV} \), what remains are only \( y' \) and \( v' \) in the mirror sector, \( g_*(T_d) = 7.25 \), while in the SM sector \( g_*(T) = 61.75 \). At \( T = T_\gamma < m_e \), \( g_*(T) = 3.9 \) while \( g_*(T') = 7.25 \), thus resulting in \( T' = 0.4 T_\gamma \). To achieve the above temperature difference, alternatively, ref. \[39\] proposed to heat up the SM sector with respect to the mirror sector by having GeV right-handed neutrinos which decay preferentially to SM particles. In sec. \[IV A\] we have taken \( T' = 0.4 T_\gamma \) which is consistent with the bound on \( \Delta N_{\text{eff}} \) (see eq. \[15\] and the discussion below).

In mirror twin Higgs models, neutrino masses can be generated from dimension five effective operators, with \( \mathbb{Z}_2 \) symmetric couplings, generated from a seesaw mechanism \[34\]. The effective Lagrangian reads

\[
L_{\text{eff}} = \frac{y_D}{M_1} (L H)^2 + \frac{y_H}{M_1} (L' H')^2 + \frac{\lambda}{M_2} (L H)(L' H') ,
\]

where \( L \) and \( L' \) are respectively the SM lepton and mirror lepton doublets. With \( M_2 \gg M_1 \), neutrinos acquire mostly Dirac masses after \( H \) and \( H' \) acquire vevs. Due to the \( \mathbb{Z}_2 \) symmetry, \( m_{\nu'} \) and \( m_\nu \) differ only by \( v'/v \),

\[
m_{\nu'} = \left( \frac{v'}{v} \right)^2 m_\nu .
\]

If on the contrary \( M_1 \gg M_2 \), neutrinos will be Majorana with their masses generated by the last term in \[22\] and so

\[
m_{\nu'} = m_\nu .
\]

Cases \[23\] and \[24\] correspond to the two mass spectra we have considered in our analysis.

As a final remark, in our scenario, the mirror photon should acquire a small mass of the order of \( 10^{-10} \text{eV} \). The broken mirror QED can be achieved by having a soft \( \mathbb{Z}_2 \) breaking mass for the mirror hypercharge gauge boson \[32\]. Furthermore, the required small kinetic mixing \( \epsilon \lesssim 10^{-5} \) in our scenario also implies the existence of millicharged (mirror) particles where the current constraints on them with mass \( \gtrsim 10^{-2} \text{GeV} \) is rather loose \( \epsilon \lesssim 10^{-4} \[40\].

VI. CONCLUSIONS

In this letter we have entertained the possibility that the anomalous spectral feature recently reported by EDGES arises from extra radiation (soft photons) injected in the bath at early times. We first have considered SM neutrino decays induced by neutrino magnetic and electric transition moments \( \mu_{\text{eff}}, V_i \rightarrow v_j + \gamma \). Treating \( \mu_{\text{eff}} \) as a free parameter and assuming full Lyman-alpha coupling, we calculated the photon flux required to address the EDGES signal. We find that an explanation based on electromagnetic-induced neutrino decays requires values for \( \mu_{\text{eff}} \) already ruled out by stellar cooling considerations, thus ruling out such possibility.

We have shown that in the presence of mirror neutrinos (as expected e.g. in mirror twin Higgs models), a larger mirror transition moment allows a sufficiently large photon flux that can account for the EDGES signal, while simultaneously satisfying astrophysical and laboratory bounds on \( \mu_{\text{eff}} \) and kinetic mixing \( \epsilon \). The mechanism generates the appropriate amount of radiation in a two-step process. In a first stage mirror neutrino decays \( v_i' \rightarrow v_j' + \gamma' \) populate the bath with a dark photon density. The decays can occur way before recombination provided they happen after the SM and mirror sectors have decoupled. In a second stage, the dark photon population is processed into visible radiation through resonant \( \gamma' - \gamma \) conversion. The resulting additional photon density thus arises from the mirror neutrino decays (after conversion) and the subdominant mirror neutrino background (CMB\(^3\)). In contrast to the first stage, the second should take place near or after recombination, to avoid the resulting photons from being totally absorbed by the medium. Assuming as well full Lyman-alpha coupling we have explicitly shown that this mechanism can raise the radiation temperature at the levels required by the EDGES anomalous spectral feature.

Finally, we showed that the scenarios we considered can be realized naturally in mirror twin Higgs models. Thus, on top of the phenomenology which come along with them, we have provided another avenue to probe them during the cosmic dark ages.

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