Type-II Seesaw at Collider, Lepton Asymmetry and Singlet Scalar Dark Matter

John McDonald* and Narendra Sahu†
Cosmology and Astroparticle Physics Group, University of Lancaster, Lancaster LA1 4YB, UK

Utpal Sarkar‡
Theory Division, Physical Research Laboratory, Navarangpura, Ahmedabad, 380 009, India

We propose an extension of the standard model with a B-L global symmetry that is broken softly at the TeV scale. The neutrinos acquire masses through a type-II seesaw while the lepton (L) asymmetry arises in the *singlet sector* but without B-L violation. The model has the virtue that the scale of L-number violation (Λ) giving rise to neutrino masses is independent of the scale of leptogenesis (Λ'). As a result the model can explain *neutrino masses, singlet scalar dark matter and leptogenesis at the TeV scale*. The stability of the dark matter is ensured by a surviving Z_2 symmetry, which could be lifted at the Planck scale and thereby allowing Planck scale-suppressed decay of singlet scalar dark matter particles of mass ≈ 3 MeV to e^+e^- pairs in the Galactic halo. The model also predicts a few hundred GeV doubly charged scalar and a long lived charged fermion, whose decay can be studied at Large Hadron Collider (LHC) and International Linear Collider (ILC).

PACS numbers: 12.60.Fr, 14.60.St, 95.35.+d, 98.80.Cq

I. INTRODUCTION

Within the standard model (SM), neutrinos are massless. On the other hand, the current low energy neutrino oscillation data [1, 2, 3] indicate that at least two of the physical left-handed (LH) neutrinos have tiny masses and therefore mix among themselves. However, as yet we do not know if the neutrinos are Dirac or Majorana. If the neutrinos are assumed to be Majorana then the sub-eV neutrino masses can be generated through the dimension five operator [4]

$$O_\nu = \frac{\phi\phi\ell\ell}{\Lambda}\,, \tag{1}$$

where Λ is the scale of lepton (L) number violation ($\Delta L = 2$). The dimension five operator (1) can originate through the celebrated see-saw mechanism.

In type-I seesaw models, three right-handed (RH) neutrinos (N's) are added without extending the gauge group of the SM. The canonical seesaw (or type-I seesaw) [5] then gives the light neutrino mass matrix:

$$m_\nu = m_\nu^I = -m_D M_N^{-1} m_D^T\,, \tag{2}$$

where m_D is the Dirac mass matrix of the neutrinos connecting the LH neutrinos (ν_L) with the RH neutrinos and M_N is the Majorana mass matrix of the RH heavy neutrinos, which also sets the scale of L-number violation (Λ). The Dirac mass terms determine the L-numbers of the RH neutrinos to be $+1$ and hence the Majorana mass of the RH neutrinos violates L-number by two units. The decays of the RH neutrinos would then violate L-number and their CP violating out-of-equilibrium decay to SM fields can be a natural source of L-asymmetry [6] in the early Universe. The CP-violation, which comes from the Yukawa couplings that determine the Dirac mass mass of neutrinos via the one-loop radiative vertex correction, requires at least two RH neutrinos. The masses of the RH neutrinos producing the final L-asymmetry then satisfy [7] [1].

$$M_N \geq O(10^9)GeV\,. \tag{3}$$

If the corresponding theory of matter is supersymmetric (SUSY) then this bound, being dangerously close to the maximum reheat temperature, poses a problem. A modest solution was proposed in ref. [9] by introducing an extra singlet heavy fermion. However, the model only achieves a reduction of above bound [7] by an order of magnitude.

In the type-II seesaw models, on the other hand, $SU(2)_L$ triplet Higgses (Δ's) are added to the SM gauge group. Explicit breaking of the L-number by trilinear couplings of the triplet Higgs scalar then induces a tiny vacuum expectation value (VEV) of the heavy triplet Higgs scalars [10], generating a light neutrino mass matrix:

$$m_\nu = m_\nu^{II} = f\mu\frac{v^2}{M_\Delta^2}\,, \tag{4}$$

where M_Δ is the mass of the triplet Higgs scalar Δ, μ is the coupling constant with mass dimension 1 for the trilinear term with the triplet Higgs and two standard model Higgs doublets, and f is the Yukawa coupling of the triplet Higgs to the light leptons. M_Δ and μ are of the same order of magnitude and set the scale (Λ) of L-number violation. v is the VEV of the SM Higgs doublet. In these models, the L-asymmetry is generated through the L-number violating decays of the Δ to SM leptons and Higgs [11]. The CP-violation, originating from the one-loop self-energy correction, requires at least two triplets. The scale of L-number violation is determined by M_Δ and μ and is

*Electronic address: j.mcdonald@lancaster.ac.uk
†Electronic address: n.sahu@lancaster.ac.uk
‡Electronic address: utpal@prl.res.in

[1] This could be next-to-lightest RH neutrino if flavor leptogenesis is considered [8].

required to be very high. In most cases the scale of L-violation in type-II seesaw models is larger than the type-I seesaw models [12].

From the above discussion we see that the L-number violating scales $M_N \sim \Lambda$ in type-I see-saw models or $M_\Delta \sim \Lambda$ in type-II seesaw models comes out to be large because the same L-number violation gives rise to both neutrino masses and mixings and to the L-asymmetry. While sub-eV neutrino masses require large values of Λ, the large values of M_N and M_Δ can easily satisfy the out-of-equilibrium decay condition, a necessary condition for the generation an L-asymmetry, without any fine-tuning of the Yukawa couplings. From this perspective, these models are attractive. However, they cannot be verified in future colliders since the scale of L-number violation is very high. Alternatives to these models are provided by mechanisms which work at the TeV scale, either in SUSY extensions of the SM [13] or by introducing an additional source of CP violation into the model [14].

On the other hand, one could ask if the origin of neutrino masses and leptogenesis is different? It is well-known that there is no one-to-one correspondence between the parameters in the neutrino mass matrix and those involved in leptogenesis [15]. In particular, in type-I see-saw model with 3 RH neutrinos, while 15 parameters enter into leptogenesis, there are only 9 parameters: 3 masses, 3 mixing angles and 3 phases (one L-number conserving phase called the Dirac phase and two L-number violating phases called Majorana phases) in the low energy neutrino mass matrix. Obviously there is no connection. Conservatively, if one considers type-I seesaw model with 2 RH neutrinos for leptogenesis as well as for low energy sub-eV neutrino masses, then the number of parameters in both cases are 9. However, one can show that there is no one-to-one correspondence between the Majorana phases responsible for L-number violation in leptogenesis and Majorana phases in neutrino mass matrix [16]. The connection between leptogenesis and low energy neutrino mass matrix is even worse in case of type-II see-saw models.

Motivated by the fact that there is no one-to-one correspondence between leptogenesis and the effective low energy neutrino mass matrix, we propose a new mechanism of leptogenesis which occurs completely in the *singlet sector* at the TeV scale [17]. As we will show, the origin of neutrino masses is different from the origin of L-asymmetry. The L-asymmetry then arises without any B-L violation. We will show that the B-L violation required for neutrino masses does not conflict with the leptogenesis. This model is then extended to incorporate a singlet scalar dark matter particle. While the nature of dark matter is still a mystery, we will show that the singlet dark scalar can either be collisionless cold dark matter (CCDM) or self-interacting dark matter (SIDM). The stability of the dark matter is ensured by a surviving Z_2 symmetry which could be lifted at the Planck scale and thereby allow Planck scale-suppressed decay of the singlet dark scalars to e^+e^- pairs in the Galactic halo. The most important feature of the model is that it predicts a few hundred GeV doubly charged scalar and a long lived singly charged fermion whose decay can be studied at the LHC/ILC.

The paper is organized as follows. In section II we introduce a model which simultaneously explains neutrino masses, singlet scalar dark matter and singlet leptogenesis. In section III, the type-II seesaw model of neutrino masses and the viability of testing it at colliders are discussed. In section IV we give a brief description of singlet leptogenesis, which arises from a conserved B-L symmetry. Section V is devoted to singlet scalar dark matter and its Planck scale suppressed decay to e^+e^- pair in the Galactic halo. In section VI we give a brief description of collider signatures and section VII concludes.

II. THE MODEL: $\mathrm{SM} \times U(1)_{B-L}$

We extend the SM gauge symmetry with a global $U(1)_{B-L}$ symmetry, which is softly broken at the TeV scale. In addition to the quarks, leptons (ℓ) and the usual Higgs doublet ϕ of the SM, we introduce two triplet scalars ξ and Δ and a singlet complex scalar χ. We also introduce two neutral singlet heavy scalars $S_a, a = 1,2$ and a charged fermion η^-. The particle content of the model and their respective quantum numbers is given in table (I).

TABLE I: Fermions and scalars included in the model.

Particle Content	$SU(3)_C \times SU(2)_L \times U(1)_Y$	$U(1)_{B-L}$
ℓ	(1,2,-1)	-1
e_R^-	(1,1,-2)	-1
ϕ	(1,2,1)	0
ξ	(1,3,2)	2
Δ	(1,3,2)	0
χ	(1,1,0)	0
η_L^-, η_R^-	(1,1,-2)	-1
S_a	(1,1,0)	0

It is then straightforward to write down the Lagrangian invariant under the SM and the global $U(1)_{B-L}$. We present here only those terms in the Lagrangian that are directly relevant to the rest of our discussions. Those are given by

$$\begin{aligned}
-\mathcal{L} \supseteq\ & f_{ij}\xi\ell_{iL}\ell_{jL} + \mu\Delta^\dagger\phi\phi + M_\xi^2\xi^\dagger\xi + M_\Delta^2\Delta^\dagger\Delta \\
& + h_{ia}\overline{e_{iR}}S_a\eta_L^- + M_{S_{ab}}^2 S_a^\dagger S_b + y_{ij}\phi\overline{\ell_{iL}}e_{jR} \\
& + g_{i1}\overline{\ell_{iL}}\phi\eta_R^- + M_\eta\overline{\eta_L^-}\eta_R^- + \mu_{\phi a}S_a\phi^\dagger\phi \\
& + \mu_{\chi a}S_a\chi^\dagger\chi + V(\phi,\chi,S,\Delta) + h.c.,
\end{aligned} \tag{5}$$

where $V(\phi,\chi,S,\Delta)$ constitutes all possible quadratic and quartic terms invariant under the SM and the global $U(1)_{B-L}$,

$$\begin{aligned}
V(\phi,\chi,S,\Delta) \ =\ & m_1^2\phi^\dagger\phi + m_2^2\chi^\dagger\chi + \lambda_1(\phi^\dagger\phi)^2 + \lambda_2(\chi^\dagger\chi)^2 \\
+\lambda_3(\phi^\dagger\phi)(\chi^\dagger\chi) \ +\ & \lambda_a(S_a^\dagger S_a)(\phi^\dagger\phi) + \lambda_b(S_b^\dagger S_b)(\chi^\dagger\chi) \\
+\lambda_4(\Delta^\dagger\Delta)(\phi^\dagger\phi) \ +\ & \lambda_5(\Delta^\dagger\Delta)(\chi^\dagger\chi).
\end{aligned} \tag{6}$$

As the Universe expands, Δ acquires a very small VEV,

$$\langle\Delta\rangle = -\mu\frac{v^2}{M_\Delta^2}\,. \quad (7)$$

For $\mu \sim M_\Delta \sim 10^{12}$ GeV and $v = \langle\phi\rangle = 174$ GeV one can get the VEV of Δ to be a few eV. The singlet scalars S_1 and S_2 acquire VEV much below the mass scale of Δ. They develop VEV at a temperature $T \sim 10$ TeV. The VEVs of S_1 and S_2 are:

$$\langle S_a\rangle = -\mu_{\phi a}\frac{v^2}{M_{S_a}^2}\,,\ a = 1,2\,. \quad (8)$$

For $\mu_{\phi a} \sim M_{S_a} \sim 10$ TeV, the VEV of S_a are in the order of 1 GeV. The singlet scalar χ does not acquire a VEV. With a Z_2 symmetry under which χ is odd while all other fields are even, χ becomes a candidate for dark matter [18, 19, 20]. As we will see below, the L-asymmetry arises from the L-number conserving decay of S_a to e_R^- and η_L^+ while the neutrino masses arise from the L-violating interactions produced via a soft breaking interaction of ξ and Δ.

III. SOFT BREAKING OF B-L SYMMETRY AND NEUTRINO MASSES

The VEV of Δ does not break B-L gauge symmetry since it is inert under the $U(1)_{\rm B-L}$ symmetry. This also ensures that B-L is an exact symmetry until it is broken softly at the TeV scale. Note that ξ does not acquire any VEV at the tree level and hence there are no neutrino masses unless the B-L symmetry is broken. We assume that $M_\xi \ll M_\Delta$ and that ξ and Δ contribute equally to the effective neutrino masses.

To generate neutrino masses we need to break the global $U(1)_{\rm B-L}$ symmetry without destroying the renormalizability of the theory while ensuring that there is no massless Nambu-Goldstone boson that can cause conflict with phenomenology. This can be achieved by adding a soft term in the Lagrangian (5)

$$-\mathcal{L}_{\Delta\xi} = m_s^2\Delta^\dagger\xi + h.c.\,, \quad (9)$$

where the mass parameter m_s is of the order of a few hundred GeV. We assume that such a soft term could originate from theories with larger symmetries. However, we will not consider its origin in this paper. The mixing between ξ and Δ is then parameterized by

$$\tan 2\theta = \frac{2m_s^2}{M_\Delta^2 - M_\xi^2} \quad (10)$$

Since we have assumed that $M_\Delta \gg M_\xi$, the mixing angle is simply

$$\theta \simeq \frac{m_s^2}{M_\Delta^2}\,. \quad (11)$$

As a result the mass eigenstates are:

$$\xi' = \xi - \left(\frac{m_s^2}{M_\Delta^2}\right)\Delta \simeq \xi \text{ and } \Delta' = \Delta + \left(\frac{m_s^2}{M_\Delta^2}\right)\xi \simeq \Delta\,. \quad (12)$$

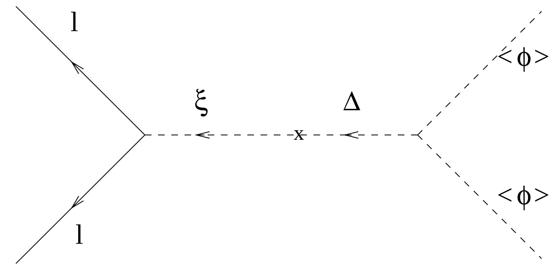

FIG. 1: Modified type-II seesaw for neutrino masses arising from soft L-number violation.

Since the soft term (9) introduces L-number violation by two units, the neutrino can acquire a mass. The effective L-number violating Lagrangian is [2]:

$$-\mathcal{L}_{\nu-mass} = f_{ij}\xi\ell_{iL}\ell_{jL} + \mu\frac{m_s^2}{M_\Delta^2}\xi^\dagger\phi\phi + f_{ij}\frac{m_s^2}{M_\Delta^2}\Delta\ell_{iL}\ell_{jL} + \mu\Delta^\dagger\phi\phi + h.c.\,. \quad (13)$$

After electroweak symmetry breaking, the field ξ acquires an induced VEV,

$$\langle\xi\rangle = -\mu\frac{v^2 m_s^2}{M_\xi^2 M_\Delta^2}\,. \quad (14)$$

This can be verified by minimization of the complete potential. The VEVs of ξ and Δ will contribute equally to neutrino masses and thus the neutrino mass matrix, derived from fig. (1), is given by

$$(m_\nu)_{ij} = -f_{ij}\mu\frac{v^2 m_s^2}{M_\xi^2 M_\Delta^2}\,. \quad (15)$$

If we consider the mass scales $\mu \sim M_\Delta \sim 10^{12}$ GeV, $m_s \sim 100$ GeV and $M_\xi \sim v$, a natural choice of the Majorana Yukawa coupling f gives the scale of neutrino masses to be $m_\nu \sim \mathcal{O}(1)$ eV, as required by laboratory, atmospheric and solar neutrino experiments. As discussed previously, one of the triplet Higgs scalar ξ could remain very light without any conflict with neutrino masses or any other phenomenology. Since the mass of ξ could be in the range of a few hundred GeV, its decay through same sign dilepton can be tested at the LHC or ILC. Thus the proposed type-II seesaw is testable in contrast to the conventional type-II seesaw. We will come back to this point in section VI while discussing collider signatures.

IV. LEPTON ASYMMETRY FROM CONSERVED B-L

We note that the interaction $S_a\eta_L^- e_R^+$ conserves B-L-number. Therefore, out-of-equilibrium decay of S_a cannot

[2] We thank Hiroaki Sugiyama for pointing out a typographical mistake in equation (13).

generate any B-L asymmetry. However, if there is CP-violation in the decay of S_a then it can produce an equal and opposite B-L asymmetry between $e_R^+(e_R^-)$ and $\eta_L^-(\eta_L^+)$. If these two B-L asymmetries never equilibrate with each other before the electroweak phase transition then the B-L asymmetry in e_R^- can be transferred to the left-handed fields in SM, while keeping an equal and opposite B-L asymmetry in η_L^+. The Yukawa couplings of the charged leptons with the SM Higgs field will transfer any asymmetry in e_R^- into an asymmetry in e_L^-, which will then take part in sphaleron processes before the electroweak phase transition. Thus the B-L asymmetry in e_R^- will generate the required baryon asymmetry via sphaleron transitions, while the equal and opposite amount of B-L asymmetry in η_L^+ will remain unaffected even after the electroweak phase transition. Eventually this field will decay slowly after the electroweak phase transition, which can generate a small L-asymmetry, but the baryon asymmetry of the universe will not be affected at this time since the sphaleron processes are out of thermal equilibrium. Therefore the baryon asymmetry will survive.

As the Universe expands, the temperature of the thermal bath falls. Below their mass scales the CP violating decay of the heavy singlet scalars $S_a, a = 1,2$ generate an equal and opposite B-L-asymmetry between $e_R^-(e_R^+)$ and $\eta_L^+(\eta_L^-)$ through

$$\begin{aligned} S_a &\rightarrow e_{iR}^- + \eta_L^+ \\ &\rightarrow e_{iR}^+ + \eta_L^- \,. \end{aligned}$$

The decay rate can be given as

$$\Gamma_a = \frac{(h^\dagger h)_{aa}}{16\pi} M_{S_a}\,, \tag{16}$$

where the masses of η_L^- and e_R^- are small in comparison to the mass of S_a. A net CP asymmetry is generated through the interference of tree-level diagram and the one-loop self-energy correction diagram involving ϕ and χ as shown in figure (2). The CP-asymmetry in the decay of S_a can be estimated as

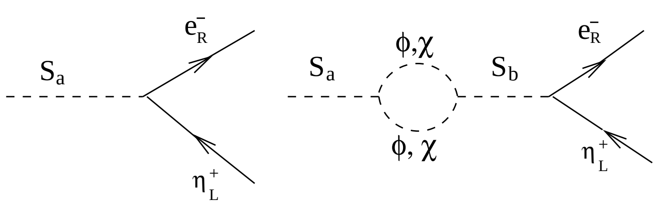

FIG. 2: The tree level diagram and the self energy correction diagram of S_a which give a net CP violation.

$$\varepsilon_a = \frac{\mathrm{Im}\left[(\mu_{\phi 1}\mu_{\phi 2}^* + \mu_{\chi 1}\mu_{\chi 2}^*)\sum_i h_{i1}h_{i2}^*\right]}{8\pi^2(M_{S_1}^2 - M_{S_2}^2)}\left[\frac{M_{S_a}}{\Gamma_a}\right]. \tag{17}$$

The lightest of S_a will generate an equal amount of e_R and η asymmetries. If the masses of S_1 and S_2 are close enough then the CP asymmetry can be resonantly enhanced, [21, 22] such that the mass scale of S_a can be a few TeV.

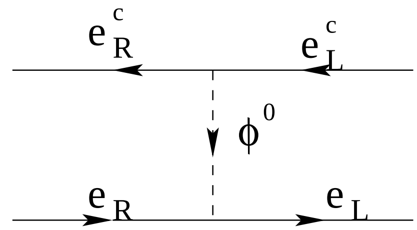

FIG. 3: The L-number conserving process which transfer the B-L asymmetry from right handed sector to the left-handed sector.

The L-asymmetry in e_R can be transferred to e_L through the t-channel process $e_R e_R^c \leftrightarrow \phi_0 \leftrightarrow e_L e_L^c$ as shown in the figure (3). These interactions will be in equilibrium for all the three generations of charged leptons below 10^5 GeV and hence there will be equal amount of e_R and e_L asymmetry. This B-L asymmetry in e_L will be converted to the baryon asymmetry of the universe before the electroweak phase transition when the sphaleron processes are in thermal equilibrium. An equal and opposite amount of B-L asymmetry in η_L will remain unaffected by these interactions.

Note that the generated B-L asymmetry in the left-handed sector is not washed out by the L-violating interactions mediated by ξ and Δ because those processes are suppressed. In particular $\ell\ell \leftrightarrow \phi\phi$ is suppressed by m_s^2/M_Δ^2. For $m_s \sim 1$ 100 GeV and $M_\Delta \sim 10^{12}$ GeV the suppression of $\Delta L = 2$ processes are of the order 10^{-20}. However, this asymmetry can be washed out through the decay: $\eta_R^+ \rightarrow \ell + \phi$, unless the decay rate

$$\Gamma_\eta = \frac{|g_{i1}|^2}{16\pi} M_\eta\,, \tag{18}$$

satisfies $\Gamma_\eta^{-1} \equiv \tau_\eta > \tau_{EW}$, where $\tau_{EW} \sim 10^{-12}$s is the time of electroweak phase transition. Furthermore, η^+ should be decayed away well before Big-Bang Nucleosynthesis (BBN) in order not to conflict with the prediction of BBN. Therefore

$$\tau_\eta < \tau_{BB} \sim 1s\,. \tag{19}$$

From Eqns. (18) and (19), and using $\tau \approx H^{-1} \approx (T^2/M_{Pl})^{-1}$, we get

$$T_{BB}^2/M_{Pl} \lesssim \frac{|g_{i1}|^2}{16\pi} M_\eta \lesssim T_{EW}^2/M_{Pl}\,, \tag{20}$$

where T_{EW} and T_{BB} respectively are the temperatures corresponding to the electroweak phase transition and BBN. We show the allowed masses and Yukawa couplings of η as a plot of $\ln|g|$ versus M_η. From figure (4) it can be seen that g can vary from 10^{-7} to 10^{-12} for M_η taking its values between 200 GeV to 1 TeV.

V. SINGLET SCALAR DARK MATTER

As the universe expands the temperature of the thermal bath falls. As a result the heavy fields S_a and Δ acquires VEV

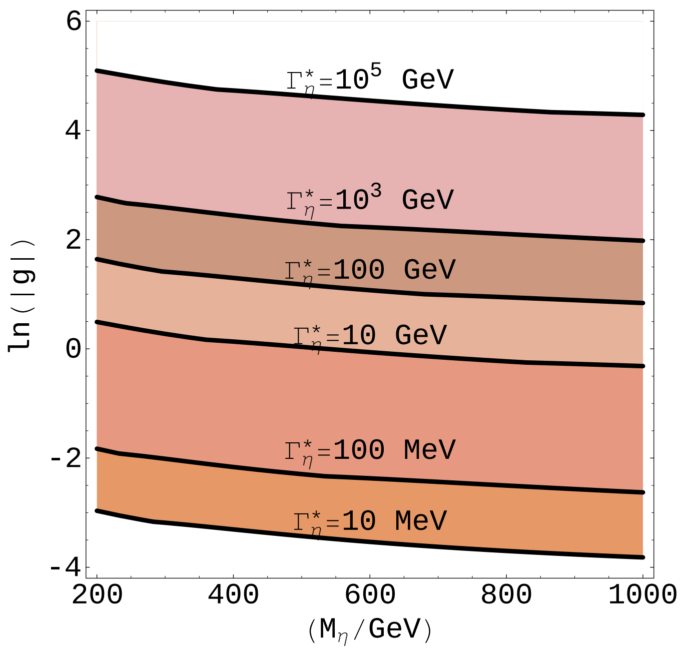

FIG. 4: Allowed contours of $\Gamma^*_\eta \equiv 10^{19}\Gamma_\eta$, required for generating a successful L-asymmetry, are shown in the plane of $\ln|g|$ versus M_η.

below their mass scales. Consequently the effective potential before the electroweak phase transition is given by

$$\begin{aligned} V_{eff} &= m_\phi^2\phi^\dagger\phi + \lambda_1(\phi^\dagger\phi)^2 + m_\chi^2\chi^\dagger\chi \\ &+ \lambda_2(\chi^\dagger\chi)^2 + \lambda_3(\phi^\dagger\phi)(\chi^\dagger\chi) \end{aligned} \quad (21)$$

where

$$\begin{aligned} m_\phi^2 &= \left(m_1^2 + \lambda_a\langle S_a\rangle^2 + \lambda_4\langle\Delta\rangle^2\right) \\ \text{and} \quad m_\chi^2 &= \left(m_2^2 + \lambda_b\langle S_b\rangle^2 + \lambda_5\langle\Delta\rangle^2\right). \end{aligned} \quad (22)$$

The above effective potential is bounded from below if and only if $\lambda_1, \lambda_2 > 0$ and $\lambda_3 > -2\sqrt{\lambda_1\lambda_2}$. For $m_\phi^2 < 0$ and $m_\chi^2 > 0$, the minimum of the potential is given by

$$\langle\phi\rangle = \begin{pmatrix} 0 \\ v \end{pmatrix} \quad \text{and} \quad \langle\chi\rangle = 0. \quad (23)$$

The VEV of ϕ gives masses to the SM fermions and gauge bosons. The physical mass of the SM Higgs is then given by $m_h = \sqrt{4\lambda_1 v^2}$. Since χ is odd under the surviving Z_2 symmetry it cannot decay to any of the conventional SM fields and hence the χ can constitute the dark matter component of the Universe.

A. Cold Dark Matter

Gauge singlet scalars interacting via the renormalisable $\chi^\dagger\chi\phi^\dagger\phi$ coupling to the Higgs doublet can account for cold dark matter (CDM) [18, 19, 20][3]. For $\lambda_3 \gtrsim 0.01$, the gauge singlet scalar CDM density is produced via conventional freeze-out from thermal equilibrium. For 10GeV $\lesssim m_{\chi\,tot} \lesssim$ 1TeV and $0.01 \lesssim \lambda_3 \lesssim 1$, the scalar density is naturally of the order of the observed CDM density over a wide region of the $(\lambda_3, m_{\chi\,tot})$ parameter space [18]. ($m_{\chi\,tot}$ denotes the physical χ mass, including the contribution from the Higgs expectation value, $m_{\chi\,tot}^2 = m_\chi^2 + \lambda_3 < \phi^\dagger\phi >$.) Such scalars have annihilation and nuclear scattering cross-sections which make them a potentially detectable CDM candidate in cryogenic detectors and neutrino telescopes, on a par with more conventional WIMP candidates [18]. They may also be produced at the LHC via Higgs decay [18, 20].

For small λ_3 there is an alternative possibility for gauge singlet scalar dark matter [19]. If $\lambda_3 \ll 1$, gauge singlet scalars in thermal equilibrium are unable to annihilate efficiently, resulting in too much CDM after freeze-out [23]. In order to evade this problem, λ_3 must be sufficiently small that the gauge singlets never come into thermal equilibrium. Although there is no freeze-out thermal relic density in this case, it is still possible to produce a CDM density via decay of thermal equilibrium Standard Model particles to χ pairs, in particular Higgs decay [19]. In this case the density of CDM is given by [19]

$$\Omega_\chi = 0.3\left(\frac{\lambda_3}{2\times10^{-10}}\right)^2\left(\frac{115\text{GeV}}{m_h}\right)^3\left(\frac{0.7}{h}\right)^2\left(\frac{m_{\chi\,tot}}{4.8\text{MeV}}\right)^2, \quad (24)$$

where m_h is the physical Higgs boson mass and h parameterizes the Hubble constant. Thus $\lambda_3 \approx 10^{-10}$ is necessary to account for CDM. A particularly interesting possibility is that χ gains its mass mostly from the Higgs expectation value. In general $m_{\chi\,tot}^2 = m_\chi^2 + \lambda_3 v^2/2$ ($v = 174$ GeV). Therefore for $m_\chi = 0$, the χ mass is a function only of λ_3. In this case the dark matter density is entirely determined by $m_{\chi\,tot}$ and m_h [19],

$$\Omega_\chi = 0.3\left(\frac{115\text{GeV}}{m_h}\right)^3\left(\frac{0.7}{h}\right)^2\left(\frac{m_{\chi\,tot}}{2.8\text{MeV}}\right)^5. \quad (25)$$

Thus $m_{\chi\,tot} \approx 3$ MeV is predicted for the scalar mass in this case. (Note that this is a rather precise estimate, due to the high power of $m_{\chi\,tot}$ in Eq.(25).) As discussed below, this is a physically significant mass scale. Light gauge singlet scalars with masses in the 1-10 MeV range [19, 24] can account for self-interacting dark matter (SIDM) [25], whilst the decay of 1-6 MeV scalars to e^+e^- pairs can account for the 511 keV γ line from the galactic center [26, 27, 28] observed by INTEGRAL [29].

B. Self-Interacting Dark Matter

SIDM has been suggested as a way to explain the absence of peaked galaxy halo profiles and the small number

[3] The notation used in [18, 19] is such that $S \leftrightarrow \chi$, $\lambda_S \leftrightarrow \lambda_3$, $m \leftrightarrow m_\chi$, $m_S \leftrightarrow m_{\chi\,tot}$ and $\eta \leftrightarrow 4\lambda_2$.

of sub-halos in galaxies as compared with CDM simulations [25]. This requires that the self-interaction cross-section of the dark matter, σ, satisfies $r_\chi \equiv \sigma/m_\chi = (0.45-5.6)\mathrm{cm}^2\mathrm{g}^{-1}$ [25]. Recently, comparison of the theoretical dynamics of the merging galaxy cluster 1E 0657-56 (the 'Bullet cluster') with observation has placed a new upper bound on this range, $r_\chi \lesssim 1\mathrm{cm}^2\mathrm{g}^{-1}$ [30]. Nevertheless, a window still remains where SIDM scalars could have a significant effect on galaxy halos.

For a gauge singlet self-interaction $\lambda_2|\chi|^4$, the total scattering cross-section is [19]

$$\sigma = \sigma_{\chi\chi^\dagger\to\chi\chi^\dagger} + \sigma_{\chi\chi\to\chi\chi} = \frac{3\lambda_2^2}{8\pi m_\chi^2}\ . \tag{26}$$

This self-interaction can modify galaxy halos if the mass is in the range $m_S = \alpha_{\lambda_2}^{1/3}\,(38.8-90.1)$ MeV (with $\alpha_{\lambda_2} = \lambda_2^2/4\pi$), where the lower bound is the value at which galaxy halos would evaporate due to interaction with hot particles in cluster halos, while the upper bound is the value at which the scalars do not interact within a typical galactic halo during a Hubble time and so have no effect [19, 25]. (The new constraint from the Bullet cluster raises the lower bound to $68.7\alpha_{\lambda_2}^{1/3}$ MeV.) If we consider the magnitude of the χ self-coupling to be similar to the SM Higgs self-coupling, such that $\lambda_2 \approx 0.025$, then this mass range is $m_\chi \sim 1-4$ MeV. In this case there is a good possibility that light scalar dark matter will have an observable effect on galaxy halos.

C. Decaying χ Dark Matter and the INTEGRAL 511 keV flux

In order to explain the 511 keV γ-ray flux observed by INTEGRAL, we need a large number density of dark matter particles at the center of the galaxy plus low energy positrons from their decay in order that the positrons can slow to non-relativistic velocities before annihilating with background electrons. These conditions require that the decaying particles are in the range 1-6 MeV, corresponding to positron injection energies $\lesssim 3$ MeV [28], in good agreement with $m_\chi \approx 3$ MeV. For an interaction of the form $\chi m_e \bar{e} e/M_*$, the γ flux relative to the experimentally observed flux is [27]

$$\frac{\Phi}{\Phi_{exp}} \approx 25\left(\frac{10^{19}\mathrm{GeV}}{M_*}\right)^2\left(\frac{\Omega_\chi}{0.23}\right)\ , \tag{27}$$

where the dominant decay mode is assumed to be $\chi \to e^+e^-$ with decay rate

$$\Gamma_{\chi\to e^+e^-} = \frac{m_e^2 m_\chi}{8\pi M_*^2}\ . \tag{28}$$

Thus with a Planck-suppressed decay rate, light gauge singlet scalars of mass $m_\chi \approx 3$ MeV can account for the observed 511 keV flux.

The above decay: $\chi \to e^+e^-$ can be realized if we allow Z_2 to be violated at the Planck scale. We have then the allowed Planck scale suppressed Lagrangian

$$\begin{aligned}\mathcal{L} \;=\; & y\frac{\chi\phi\bar{\ell}e_R}{M_{\rm Pl}} + f\frac{\chi\xi\ell\ell}{M_{\rm Pl}} + f\frac{\chi(m_s^2/M_\xi^2)\Delta\ell\ell}{M_{\rm Pl}} \\ & + g\frac{\chi\bar{\ell}\phi\eta_R^-}{M_{\rm Pl}} + h\frac{\chi S\eta_L^- e_R^+}{M_{\rm Pl}} + \hat{h}\chi\eta_L^- e_R^+ + h.c.\,.\end{aligned} \tag{29}$$

Since $m_\chi \ll M_\eta$, the decay modes: $\chi \to \eta_L^- e_R^+$ and $\chi \to \bar{\ell}\phi\eta_R^-$, $S\eta_L^- e_R^+$ are forbidden. Thus the relevant effective Z_2 violating Lagrangian is given as:

$$\mathcal{L} = \frac{m_e}{M_{\rm Pl}}\chi\bar{e_L}e_R + f'\chi\ell\ell + f'(m_s^2/M_\xi^2)\chi\ell\ell + h.c.\,, \tag{30}$$

where $f' \equiv f(\langle\Delta\rangle/M_{\rm Pl}) \leq 10^{-26}$. Therefore, the contribution of the second and third terms are negligible to the observed γ-ray flux while the first term is in the right ballpark.

VI. SIGNATURES OF $\xi^{\pm\pm}$ AND $\eta^\pm$

The doubly charged component of the light triplet Higgs ξ can be observed through its decay into same sign dileptons [31]. Since $M_\Delta \gg M_\xi$, the production of Δ particles in comparison to ξ is highly suppressed. Hence it is worth looking for the signature of $\xi^{\pm\pm}$ either at LHC or ILC. From Eq. (13) one can see that the decay $\xi^{\pm\pm} \to \phi^\pm\phi^\pm$ are suppressed since the decay rate involves the factor $\frac{m_s^2}{M_\Delta^2} \sim 10^{-20}$. While the decay mode $\xi^{\pm\pm} \to h^\pm W^\pm$ is phase space suppressed, the decay mode $\xi^{\pm\pm} \to W^\pm W^\pm$ is suppressed because the VEV of ξ is small as required for sub-eV neutrino masses and to maintain the ρ parameter of SM to be unity. Therefore, once produced, ξ mostly decays through same sign dileptons: $\xi^{\pm\pm} \to \ell^\pm\ell^\pm$. Note that the doubly charged particles cannot couple to quarks. Therefore the SM background of the process $\xi^{\pm\pm} \to \ell^\pm\ell^\pm$ is quite clean and so the detection will be unmistakable. From Eq. (13) the decay rate of the process $\xi^{\pm\pm} \to \ell^\pm\ell^\pm$ is given by

$$\Gamma_{ii} = \frac{|f_{ii}|^2}{8\pi}M_{\xi^{++}} \quad and \quad \Gamma_{ij} = \frac{|f_{ij}|^2}{4\pi}M_{\xi^{++}}\,, \tag{31}$$

where the f_{ij} are highly constrained by lepton flavor violating decays. Therefore if a doubly charged scalar can be detected in the future, the neutrino mass patterns can be probed at a collider [32] [4].

Since the mass of η^- is a few hundred GeV, the decay $\eta^- \to he^-$ can be observed in future colliders (LHC/ILC). On the other hand, η^- can also be produced in high energy neutrino collisions with matter through $\overline{\nu_L}e_R^- \to \overline{\nu_L}\eta_R^-$. Since it is long-lived, it can travel large distances and be detected in neutrino telescopes [34] [5].

---

[4] There has been recent interest in detecting neutrino mass parameters at colliders [33]. These references appeared on arXiv after our paper.

[5] We thank John F. Beacom for bringing this to our notice.

VII. CONCLUSIONS

We have introduced a new leptogenesis mechanism in a $U(1)_{B-L}$ extension of the SM which allowed us to explain simultaneously neutrino masses, dark matter and leptogenesis. The important message is that the L-asymmetry arises without any B-L violation. Since the L-number violation required for leptogenesis and neutrino masses are different, the leptogenesis scale can be lowered to as low as a few TeV.

Neutrino masses arise through a modified type-II seesaw which predicts a few hundred GeV triplet scalar. Note that in conventional type-II seesaw models the mass scale of the triplets is required to be at least $O(10^{10})$ GeV to produce sub-eV neutrino masses. Since one of the triplets in our model, namely ξ, has a mass of a few hundred GeV, the proposed model can be tested in the near future at colliders through same sign dilepton decay of ξ.

The model also predicts a singly charged fermion η^- of mass ranging from 200 GeV to 1 TeV. η^- can be produced in high energy neutrino collision with matter. Since it is long-lived, it can travel large distances and be observed in neutrino telescopes.

We proposed a singlet scalar, χ, with mass ≈ 3 MeV as a candidate for dark matter, whose stability is ensured by a Z_2 symmetry. This Z_2 discrete symmetry can be broken at the Planck scale. We then showed that a possible origin of 511 KeV Galactic line detected by INTEGRAL could be the Planck scale-suppressed decay of χ to e^-e^+ pairs in the Galactic halo.

Acknowledgement

JM and NS were supported by the European Union through the Marie Curie Research and Training Network "UniverseNet" (MRTN-CT-2006-035863) and by STFC (PPARC) Grant PP/D000394/1.

[1] Q.R. Ahmed *et al* (SNO Collaboration), Phys. Rev. Lett. **89**, 011301-011302 (2002); J.N. Bahcall and C. Pena-Garay, [arXiv:hep-ph/0404061].

[2] S. Fukuda *et al* (Super-Kamiokande Collaboration), Phys. Rev. Lett. **86**, 5656 (2001).

[3] K. Eguchi *et al* (KamLAND collaboration), Phys. Rev. Lett. **90**, 021802 (2003).

[4] S. Weinberg, Phys. Rev. Lett. **43**, 1566 (1979).

[5] P. Minkowski, Phys. Lett. **B 67**, 421 (1977); M. Gell-Mann, P. Ramond and R. Slansky in *Supergravity* (P. van Niewenhuizen and D. Freedman, eds), (Amsterdam), North Holland, 1979; T. Yanagida in *Workshop on Unified Theory and Baryon number in the Universe* (O. Sawada and A. Sugamoto, eds), (Japan), KEK 1979; R.N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980) .

[6] M. Fukugita and T. Yanagida, Phys. Lett. **B174**, 45 (1986) .

[7] S. Davidson and A. Ibarra, Phy . Lett. B**535**, 25 (2002); W. Buchmuller, P. Di Bari and M. Plumacher, Nucl. Phys. B**643**, 367 (2002) .

[8] See for example G. Engelhard, Y. Grossman, E. Nardi and Y. Nir, Phys. Rev. Lett. **99** (2007) 081802 .

[9] E. Ma, N. Sahu and U. Sarkar, J. Phys. G **34**, 741 (2007) [arXiv:hep-ph/0611257]; J. Phys. G **32**, L65 (2006) [arXiv:hep-ph/0603043].

[10] J. Schechter and J.W.F. Valle, Phys. Rev. **D 22**, 2227 (1980); M. Magg and C. Wetterich, Phys. Lett. B **94**, 61 (1980); R. N. Mohapatra and G. Senjanovic, Phys. Rev. D**23**, 165 (1981); G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. B**181**, 287 (1981) .

[11] E. Ma and U. Sarkar, Phys. Rev. Lett. **80** (1998) 5716-5719 .

[12] S. Antusch and S. F. King, Phys. Lett. B **597**, 199 (2004); T. Hambye and G. Senjanovic, Phys. Lett. B**582**, 73 (2004); N. Sahu and U. Sarkar, Phys. Rev. D **74**, 093002 (2006); N. Sahu and S. Uma Sankar, Nucl. Phys. B **724**, 329 (2005); Phys. Rev. D **71**, 013006 (2005); Pei-hong Gu and Xiao-jun Bi, Phys. Rev. D**70**, 2004 (063511) .

[13] L. Boubekeur, T. Hambye and G. Senjanovic, Phys. Rev. Lett. **93**, 111601 (2004); [arXiv:hep-ph/0404038]; T. Hambye, J.M. Russel, S.M. West, JHEP **0407**, 070 (2004), [arXiv:hep-ph/0403183]; E.J. Chun, [arXiv: hep-ph/0508050];

[14] A partial list: N. Sahu and U.A. Yajnik, Phys. Rev. D**71**, 2005 (023507), [arXiv:hep-ph/0410075]; Phys. Lett. B **635** (2006) 11, [arXiv:hep-ph/0509285]; A. Sarkar and U. A. Yajnik, Phys. Rev. D **76**, 025001 (2007) [arXiv:hep-ph/0703142]; S. F. King and T. Yanagida, Prog. Theor. Phys. **114**, 1035 (2006) [arXiv:hep-ph/0411030]; A. Abada, H. Aissaoui and M. Losada, Nucl. Phys. B **728**, 55 (2005) [arXiv:hep-ph/0409343].

[15] A. S. Joshipura, E. A. Paschos and W. Rodejohann, JHEP **0108** (2001) 029; S. Pascoli, S. T. Petcov and W. Rodejohann, Phys. Rev. D **68**, 093007 (2003); J. R. Ellis and M. Raidal, Nucl. Phys. B **643** (2002) 229; G. C. Branco, T. Morozumi, B. M. Nobre and M. N. Rebelo, Nucl. Phys. B **617** (2001) 475 [arXiv:hep-ph/0107164].

[16] K. Bhattacharya, N. Sahu, U. Sarkar and S.K. Singh, Phy. Rev. D **74**, 093001 (2006); T. Endoh, S. Kaneko, S. K. Kang, T. Morozumi and M. Tanimoto, Phys. Rev. Lett. **89** (2002) 231601; P. H. Frampton, S. L. Glashow and T. Yanagida, Phys. Lett. B **548** (2002) 119 .

[17] M. Frigerio, T. Hambye and E. Ma, JCAP **0609**, 009 (2006); N. Sahu and U. Sarkar, Phys. Rev. D **76**, 045014 (2007) [arXiv:hep-ph/0701062].

[18] J. McDonald, Phys. Rev. D **50**, 3637 (1994) [arXiv:hep-ph/0702143].

[19] J. McDonald, Phys. Rev. Lett. **88**, 091304 (2002) [arXiv:hep-ph/0106249].

[20] V. Silveira and A. Zee, Phys. Lett. B **161**, 136 (1985); C. P. Burgess, M. Pospelov and T. ter Veldhuis, Nucl. Phys. B **619**, 709 (2001) [arXiv:hep-ph/0011335]; V. Barger, P. Langacker and G. Shaughnessy, Phys. Rev. D **75**, 055013 (2007) [arXiv:hep-ph/0611239]; H. Sung Cheon, S. K. Kang and C. S. Kim, arXiv:0710.2416 [hep-ph].

[21] M. Flanz, E.A. Paschos and U. Sarkar, Phys. Lett. **B 345**, 248 (1995); A. Pilaftsis and T. E. J. Underwood, Nucl. Phys. B **692**, 303 (2004) .

[22] G. D'Ambrosio, T. Hambye, A. Hektor, M. Raidal and

A. Rossi, Phys. Lett. B **604**, 199 (2004); E. J. Chun and S. Scopel, Phys. Lett. B **636**, 278 (2006); [arXiv:hep-ph/0609259] .

[23] M. C. Bento, O. Bertolami and R. Rosenfeld, Phys. Lett. B **518**, 276 (2001) [arXiv:hep-ph/0103340].

[24] M. C. Bento, O. Bertolami, R. Rosenfeld and L. Teodoro, Phys. Rev. D **62**, 041302 (2000) [arXiv:astro-ph/0003350].

[25] D. N. Spergel and P. J. Steinhardt, Phys. Rev. Lett. **84**, 3760 (2000) [arXiv:astro-ph/9909386].

[26] C. Boehm, D. Hooper, J. Silk, M. Casse and J. Paul, Phys. Rev. Lett. **92**, 101301 (2004) [arXiv:astro-ph/0309686].

[27] C. Picciotto and M. Pospelov, Phys. Lett. B **605**, 15 (2005) [arXiv:hep-ph/0402178].

[28] J. F. Beacom and H. Yuksel, Phys. Rev. Lett. **97**, 071102 (2006) [arXiv:astro-ph/0512411]; J. F. Beacom, N. F. Bell and G. Bertone, Phys. Rev. Lett. **94**, 171301 (2005) [arXiv:astro-ph/0409403].

[29] P. Jean *et al.*, Astron. Astrophys. **407**, L55 (2003) [arXiv:astro-ph/0309484]; J. Knodlseder *et al.*, Astron. Astrophys. **411**, L457 (2003) [arXiv:astro-ph/0309442].

[30] S. W. Randall, M. Markevitch, D. Clowe, A. H. Gonzalez and M. Bradac, arXiv:0704.0261 [astro-ph]; M. Markevitch *et al.*, Astrophys. J. **606**, 819 (2004) [arXiv:astro-ph/0309303].

[31] G. Barenboim, K. Huitu, J. Maalampi and M. Raidal, Phys. Lett. B **394**, 132 (1997); K. Huitu, J. Maalampi, A. Pietila and M. Raidal, Nucl. Phys. B **487**, 27 (1997); T. Han, H. E. Logan, B. Mukhopadhyaya and R. Srikanth, Phys. Rev. D **72**, 053007 (2005); E. Ma, M. Raidal and U. Sarkar, Nucl. Phys. B **615**, 313 (2001); C. Yue and S. Zhao, [arXiv: hep-ph/0701017].

[32] A. G. Akeroyd and M. Aoki, Phys. Rev. D **72**, 035011 (2005); E. J. Chun, K. Y. Lee and S. C. Park, Phys. Lett. B **566**, 142 (2003).

[33] A. G. Akeroyd, M. Aoki and H. Sugiyama, arXiv:0712.4019 [hep-ph]; M. Kadastik, M. Raidal and L. Rebane, arXiv:0712.3912 [hep-ph]; J. Garayoa and T. Schwetz, arXiv:0712.1453 [hep-ph].

[34] S. Ando, J. F. Beacom, S. Profumo and D. Rainwater, arXiv:0711.2908 [hep-ph] and the references therein.