A UV Complete Framework of Freeze-in Massive Particle Dark Matter

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Abstract

We propose a way to generate tiny couplings of freeze-in massive particle dark matter with the Standard Model particles dynamically by considering an extension of the electroweak gauge symmetry. The dark matter is considered to be a singlet under this extended gauge symmetry which we have assumed to be the one in a very widely studied scenario called left-right symmetric model. Several heavy particles, that can be thermally inaccessible in the early Universe due to their masses being greater than the reheat temperature after inflation, can play the role of portals between dark matter and Standard Model particles through effective one loop couplings. Due to the loop suppression, one can generate the required non-thermal dark matter couplings without any need of highly fine tuned Yukawa couplings beyond that of electron Yukawa with the Standard Model like Higgs boson. We show that generic values of Yukawa couplings as large as $10^{-4} - 10^{-3}$ can keep the dark matter out of thermal equilibrium in the early Universe and produce the correct relic abundance later through the freeze-in mechanism. Though the dark matter effective couplings are tiny as required by the freeze-in scenario, the associated rich particle sector of the model can be probed at ongoing and near future experiments. The allowed values of dark matter mass can remain in a wide range from keV to TeV order keeping the possibilities of warm and cold dark matter equally possible.

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I. INTRODUCTION

In view of several astrophysical and cosmological evidences, the existence of non-baryonic form of matter, or the so called Dark Matter (DM) in large amount in the present Universe has become an irrefutable fact. Among these evidences, the galaxy cluster observations by Fritz Zwicky [1] back in 1933, observations of galaxy rotation curves in 1970’s [2], the more recent observation of the bullet cluster [3] and results from several satellite borne cosmology experiments like WMAP [4] and Planck [5] are the most prominent ones. The precise measurements of the cosmology experiments reveal that more than 80% matter content of our Universe is in the form of this non-baryonic or DM form. The amount of DM present in the Universe is often expressed by a quantity $\Omega_{DM}h^2$, which is called the relic density of DM and it is the ratio of present mass density of DM by the critical density of the Universe. The value of DM relic density at the present epoch is $0.1172 \leq \Omega_{DM}h^2 \leq 0.1226$ at 67% C.L. [5]. Here $h = H_0/100$ is a parameter of order unity while $H_0$ being the present value of the Hubble parameter.

In spite these astrophysical and cosmological evidences of DM, the information regarding the constituents and origin of DM still remains unknown to us. One of the well motivated and most studied scenario is to assume the thermal origin of DM [6, 7], where DM particles were produced thermally at the early Universe and depending upon the mass and interaction strength, DM maintained both thermal and chemical equilibrium with the plasma upto a certain temperature of the Universe. Decoupling of DM from the thermal bath particles occurred at around a temperature $T_f$, which is known as the freeze-out temperature, where the interaction rate of DM dropped below the expansion rate of the Universe governed by the Hubble parameter $H$. Being decoupled from the rest of the plasma these DM particles becomes thermal relic whose density, after decoupling, is only affected by the expansion of the Universe. The list of criteria a particle DM candidate should fulfil rules out all the Standard Model (SM) particles from being DM candidates, leading to several beyond Standard Model (BSM) proposals in the last few decades. Most of the thermal DM candidates studied in the literature fall into a category called weakly interacting massive particle (WIMP) [8–10], which has mass in the range of few GeV to few TeV and weak scale couplings. The interesting coincidence that a DM particle having mass and couplings around the electroweak scale can give rise to the correct dark matter relic abundance is often referred to as the WIMP Miracle.
Now, if such type of particles whose interactions are of the order of electroweak inter-
actions really exist then we should expect their signatures in various DM direct detection
experiments where the recoil energies of detector nuclei scattered by DM particles are be-
ing measured. However, after decades of running, direct detection experiments are yet to
observe any DM-nucleon scattering [11–13]. The absence of dark matter signals from the
direct detection experiments have progressively lowered the exclusion curve in its mass-cross
section plane. With such high precision measurements, the WIMP-nucleon cross section will
soon overlap with the neutrino-nucleon cross section. Similar null results have been also re-
ported by other direct search experiments like the large hadron collider (LHC) giving upper
limits on DM interactions with the SM particles. A recent summary of collider searches for
DM can be found in [14]. Although such null results could indicate a very constrained region
of WIMP parameter space, they have also motivated the particle physics community to look
for beyond the thermal WIMP paradigm where the interaction scale of DM particle can be
much lower than the scale of weak interaction i.e. DM may be more feebly interacting than
the thermal WIMP paradigm.

One of the viable alternatives of WIMP paradigm, which may be a possible reason of
null results at various direct detection experiments, is to consider the non-thermal origin
of DM [15]. In this scenario, the initial number density of DM in the early Universe is
negligible and it is assumed that the interaction strength of DM with other particles in the
thermal bath is so feeble that it never reaches thermal equilibrium at any epoch in the early
Universe. In this set up, DM is mainly produced from the out of equilibrium decays of some
heavy particles in the plasma. It can also be produced from the scatterings of bath particles,
however if same couplings are involved in both decay as well as scattering processes then
the former has the dominant contribution to DM relic density over the latter one [15–17].
The production mechanism for non-thermal DM is known as freeze-in and the candidates
of non-thermal DM produced via freeze-in are often classified into a group called Freeze-in
(Feebly interacting) massive particle (FIMP). For a recent review of this DM paradigm,
please see [18]. Now, if the mother particle is in thermal equilibrium with the bath then
the maximum production of DM occurs when the temperature of the Universe $T \simeq M_0$, the
mass of mother particle. Therefore, the non-thermality criterion enforces the couplings to
be extremely tiny via the following condition $\left| \frac{\Gamma}{H} \right|_{T \simeq M_0} < 1$ [19], where $\Gamma$ is the decay width.
For the case of scattering, one has to replace $\Gamma$ by the interaction rate $n_{\text{eq}} \langle \sigma v \rangle$, $n_{\text{eq}}$ being the
equilibrium number density of mother particle. These types of freeze-in scenarios are known as IR-freeze-in [16, 18, 20–26] where DM production is dominated by the lowest possible temperature at which it can occur i.e. $T \sim M_0$, since for $T < M_0$, the number density of mother particle becomes Boltzmann suppressed. Here DM interacts with the visible sector via renormalizable operators (dimension $d \leq 4$) only. There may be a situation in IR-freeze-in, where mother particle itself is out of thermal equilibrium and in such cases, first one has to calculate the distribution function of mother particle considering its all possible production and decay modes. This distribution function is necessary to compute the non-thermal $^1$ averages of decay width and annihilation cross sections. Once we know these quantities, the Boltzmann equation for the non-thermal DM can be solved in terms of its comoving number density following the usual procedure [24, 25].

As mentioned earlier that to maintain a situation where DM remains out of thermal equilibrium, one needs extremely tiny couplings of DM with the particles in the plasma. However, theoretically, the origin of such extremely low values of couplings is in general, not obvious. One of the possible explanation of such feeble interactions is to consider DM to be connected to the visible sector via non-renormalizable higher dimensional effective operators. This results in a different type of freeze-in mechanism known as UV-freeze-in [15, 27, 28], where the comoving number density of DM is directly proportional to reheat temperature $T_{RH}$ of the Universe and thus sensitive to the early Universe cosmology. Another interesting way to generate tiny dimensionless couplings is through the recently proposed clockwork mechanism [29, 30] which has been recently explored in the context of freeze-in DM by the authors of [31, 32]. In this work, we try to explain the origin of such tiny couplings by considering a renormalizable gauge extension of the SM where effective FIMP couplings with the rest of the particles can arise at radiative level, leading to the required suppression naturally. As an illustrative example, we consider a left-right symmetric gauge extension of SM where the FIMP candidate is a gauge singlet. However, at one loop level the gauge bosons can decay into the FIMP, with several particles going inside the loop. The particles in the loop do have sizeable couplings with FIMP, but that that does not lead to thermal production of FIMP DM if the corresponding scattering rates always remain smaller than

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$^1$ Here we use the word non-thermal average to distinguish it from the thermal average where the Maxwell-Boltzmann distribution function is used [25].
the expansion rate of the Universe and their decay into FIMP are kinematically forbidden. Another way is to consider these mediator particles to be too heavy to be produced at the end of inflation (having mass more than the reheat temperature). We find that both these scenarios can lead to different predictions for dark sector. Such an exercise can be carried out for simpler gauge extension of SM as well, but we perform it here for left-right symmetric model (LRSM) which has several other motivations.

Rest of the article is organised as follows. In section II, we discuss our model followed by the details of the calculation of effective one loop vertex factors of dark matter interactions with the neutral gauge bosons in section III. In section IV we discuss the details of dark matter calculations, results and then finally conclude in section V

II. THE MODEL

The LRSM is one of most highly motivated BSM frameworks which in its generic form [33–37], not only explains the origin of parity violation in weak interactions but also explains the origin of tiny neutrino masses naturally. The gauge symmetry group and the field content of the generic LRSM can also be embedded within grand unified theory (GUT) symmetry groups like $SO(10)$ providing a non-supersymmetric route to gauge coupling unification. The right handed fermions of the SM forms doublet under a new $SU(2)_R$ group in LRSM such that the theory remains parity symmetric at high energy. This necessitates the inclusion of the right handed neutrino as a part of the right handed lepton doublet. To be more appropriate, the gauge symmetry of the Standard Model namely, $SU(3)_c \times SU(2)_L \times U(1)_Y$ is upgraded to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ such that the right handed fermions transform as doublets under $SU(2)_R$, making the theory left-right symmetric. The model also has an in-built discrete $Z_2$ symmetry or D-parity which ensures the equality of couplings in $SU(2)_{L,R}$ sectors. The effective parity violating electroweak physics at low energy arises as a result of spontaneous breaking of the $SU(2)_R \times U(1)_{B-L} \times D$ to $U(1)_Y$ of the SM.

The minimal LRSM however, does not contain a naturally stable DM candidate. One can of course realise a long lived keV right handed neutrino DM in these models. Such a scenario leading to warm dark matter scenarios has been investigated within LRSM in [38–40]. Due the presence of $SU(2)_R$ gauge interactions, such a right handed neutrino dark matter can be thermally produced in the early Universe, unlike in typical keV right handed neutrino
DM models where non-thermal origin is required [41]. On the other hand, to have WIMP DM type realisation, the minimal LRSM can be extended by additional scalar or fermion multiplets in the spirit of minimal DM scenario [42–44]. Such minimal dark matter scenario in LRSM has been studied recently by the authors of [45, 46]. In these models, the dark matter candidate is stabilised either by a $Z_2 = (-1)^{B-L}$ subgroup of the $U(1)_{B-L}$ gauge symmetry or due to an accidental symmetry at the renormalisable level due to the absence of any renormalisable operator leading to dark matter decay [47]. Some more studies on left-right dark matter also appeared in the recent works [48–52]. The possibility of right handed neutrino dark matter in a different version of LRSM where the right handed lepton doublets do not contain the usual charged leptons, was also studied in the recent works [53–55].

In this work, we intend to have a purely non-thermal DM within LRSM. We therefore consider gauge singlet fermion $N$ as our DM candidate. We introduce an additional $Z_2$ symmetry under which this new singlet fermion is odd and hence can be stable if it happens to be the lightest $Z_2$ odd particle. One can also consider scalar singlet DM, but scalars usually have quartic couplings with other scalars and it is often difficult to forbid them from symmetry arguments. We also introduce two copies of vector like fermion doublets $\psi$ and a pair of scalar doublets $H_{L,R}$ to the minimal LRSM. These additional fields play the role of generating interactions of the SM sector with the DM particle $N$ at radiative level, as we will see below.
The relevant Yukawa couplings for the Standard Model fermion masses can be written as

\[ \mathcal{L}^{SM}_Y = y_{ij} \bar{\ell}_{iL} \Phi \ell_{jR} + y'_{ij} \bar{\ell}_{iL} \tilde{\Phi} \ell_{jR} + Y_{ij} \bar{Q}_{iL} \Phi Q_{jR} + Y'_{ij} \bar{Q}_{iL} \tilde{\Phi} Q_{jR} + \frac{1}{2} (f_L)_{ij} \ell_{iL}^T C \sigma_2 \Delta_L \ell_{jL} + \frac{1}{2} (f_R)_{ij} \ell_{iR}^T C \sigma_2 \Delta_R \ell_{jR} + \text{H.c.} \]  

where \( \tilde{\Phi} = \tau_2 \Phi^* \tau_2 \), \( C \) is the charge conjugation operator and the indices \( i, j = 1, 2, 3 \) correspond to the three generations of fermions. The Yukawa couplings involving the new fermions can be written as

\[ \mathcal{L}^{new}_Y = Y_{\psi} \bar{\psi} H_L N + Y'_{\psi} \bar{\psi}' H_R N + M \bar{\psi} \psi + M' \bar{\psi}' \psi' + f_\psi \sigma (\bar{\psi} \psi - \bar{\psi}' \psi') + Y_{\phi} \bar{\phi} \Psi \psi + Y'_{\phi} \bar{\phi}' \tilde{\Phi} \Psi' \]

The details of the scalar potential is given in appendix A. At a very high energy scale, the parity odd singlet \( \sigma \) can acquire a vev to break D-parity spontaneously while the neutral component of \( \Delta_R \) acquires a non-zero vev at a later stage to break the gauge symmetry of the LRSM into that of the SM which then finally gets broken down to the \( U(1)_{\text{em}} \) of electromagnetism by the vev of the neutral component of Higgs bidoublet \( \Phi \). Thus, the symmetry breaking chain is

\[
SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times D \quad \langle \sigma \rangle \quad SU(2)_L \times SU(2)_R \times U(1)_{B-L} \quad \langle \Delta_R \rangle \quad SU(2)_L \times U(1)_Y \quad \langle \Phi \rangle \quad U(1)_{\text{em}}
\]

Denoting the vev of the neutral components of the bidoublet as \( k_{1,2} \) and that of triplet \( \Delta_R \) as \( v_R \), the gauge boson masses after spontaneous symmetry breaking can be written as

\[
M^2_{W_L} = \frac{g^2}{4} k_1^2, \quad M^2_{W_R} = \frac{g^2}{2} v_R^2.
\]
\[ M_{Z_L}^2 = \frac{g^2 k_1^2}{4 \cos^2 \theta_w} \left( 1 - \frac{\cos^2 2\theta_w k_1^2}{2 \cos^4 \theta_w v_R^2} \right), \quad M_{Z_R}^2 = \frac{g^2 v_R^2 \cos^2 \theta_w}{\cos 2\theta_w}, \]

where \( \theta_w \) is the Weinberg angle. The neutral components of the other scalar fields \( \psi, \Omega_{L,R} \) do not acquire any vev. However, the neutral component of the scalar triplet \( \Delta_L \) can acquire a tiny but non-zero induced vev after the electroweak symmetry breaking as

\[ v_L = \gamma \frac{M_{W_L}^2}{v_R}, \quad (3) \]

with \( M_{W_L} \sim 80.4 \text{ GeV} \) being the weak boson mass and \( \gamma \) is a function of various couplings in the scalar potential. The bidoublet also gives rise to non-zero \( W_L - W_R \) mixing parameterised by \( \xi \) as

\[ \tan 2\xi = \frac{2k_1 k_2}{v_R^2 - v_L^2}, \quad (4) \]

which is constrained to be \( \xi \leq 7.7 \times 10^{-4} \) [56, 57].

It should be noted that, our scenario can work even without the D-parity odd singlet scalar. In such a case, the parameters of the left and right sectors are equal until the \( SU(2)_R \times U(1)_{B-L} \) symmetry breaking scale.

### III. DECAY OF \( Z_{L,R} \) INTO FIMP

The effective vertex for the \( Z_{L,R} \) decay into two \( N \)'s is given as

\[ \mathcal{L} \supset A_{Z_{L,R}}^{NN} Z_{L,R} \overline{N} \gamma^\mu \gamma_5 N \]

\[ A_{Z_{L,R}}^{NN} = \frac{-ig_{L,R} Y_{\psi,\psi'}^{2}}{32\pi^2 M_{Z_{L,R}}^2} \left[ 2 M_{Z_{L,R}}^2 + M_{\psi,\psi'}^2 \ln \left( \frac{2 M_{Z_{L,R}}^2 - M_{Z_{L,R}}^2 + \sqrt{M_{Z_{L,R}}^2 - 4 M_{\psi,\psi'}^2}}{2 M_{\psi,\psi'}^2} \right)^2 \right] \]

\[ = \frac{-ig_{L,R} Y_{\psi,\psi'}^{2}}{32\pi^2 x_{L,R}} \left[ 2 x_{L,R} + \ln \left( \frac{1}{2} \left( 2 - x_{L,R} + \sqrt{x_{L,R}^2 - 4 x_{L,R}} \right) \right)^2 \right] \]

\[ A_{Z_{L,R}}^{NN} \rightarrow \frac{-ig_{L,R} Y_{\psi,\psi'}^{2}}{16\pi^2} \quad x_{L,R} \rightarrow \infty \]

\[ = \frac{-ig_{L,R} Y_{\psi,\psi'}^{2}}{32\pi^2} \quad x_{L,R} \rightarrow 0 \]

where, \( x_{L,R} = \frac{M_{Z_{L,R}}^2}{M_{\psi,\psi'}^2} \).

In the above expression, all the masses of the particles inside the loop are taken to be same that is, \( M_{\psi,\psi'} = M_{H_{L,R}} \). Now, if the mass of the decaying boson is smaller than that of the masses of the particles inside loop, then the form factor reaches a saturated value of \( \frac{1}{32\pi^2} \).
FIG. 1: Neutral gauge boson decay into FIMP

FIG. 2: Form factor for the one loop decay shown in figure 1.

On the other hand, if the degenerate loop particles are heavier than the decaying particle, then the form factor increases by a factor of 2 becoming $\frac{1}{16\pi^2}$. Now, for example, taking the decaying boson to be $Z_L$ having mass $M_{Z_L} = 91$ GeV and the masses of the particles inside, $\psi$ and $H_L$, to be $M_{\psi} = M_{H_L} = 1$ TeV, the loop factor comes out to be around $\frac{1}{32\pi^2}$. On the other hand if the decaying boson is $Z_R$ with mass $M_{Z_R} = 6$ TeV and keeping the mass of the internal particles $\psi'$ and $H_R$ to be 1 TeV then the loop factor comes out to be $0.006841$ close to $\frac{1}{16\pi^2}$. This behaviour of the form factor is summarised in the plot shown in Fig. 2.
IV. FIMP RELIC DENSITY CALCULATION

In this model, as mentioned earlier, the fermion $N$ is completely singlet under the left-right symmetry group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and it has an odd $\mathbb{Z}_2$ parity. However, being the lightest $\mathbb{Z}_2$-odd particle, all the $\mathbb{Z}_2$ parity conserving decay modes of $N$ are kinematically forbidden, making $N$ absolutely stable over the cosmological time scale. Therefore, we consider $N$ as a suitable dark matter candidate in this work. Apart from being the lightest $\mathbb{Z}_2$-odd particle, $N$ also fits into our desired FIMP scenario. This is due to the fact that, in the present model, all the portal interactions of $N$ via $Z_{L,R}$ and $\Phi$ with the Standard Model particles are one loop suppressed. This naturally makes $N$ very feebly interacting with the thermal bath and as a result $N$ remains out of thermal equilibrium. Although $N$ can thermalise with the plasma via t-channel scattering processes like $H_{L(R)} H_{L(R)} \rightarrow \bar{N} N$, $\bar{\psi}(t) \psi(t) \rightarrow \bar{N} N$, $\bar{\psi}(t) H_{L(R)} \rightarrow N Z_{L(R)}$, $\bar{\psi}(t) H_{L(R)}^+ \rightarrow N W_{L(R)}^+$ etc, where the two $N$ final states are suppressed by $Y_{\psi}^4$ while the scattering processes with single $N$ in the final state are proportional to $Y_{\psi}^2$. Since, the Yukawa couplings are $Y_{\psi} Y_{\psi'} < 1$ in general, the dominant contribution arises from the scattering processes with single $N$ in the final state. In the left panel of Fig. 3 we demonstrate the variation of ratio between the interaction rate $n_{\text{eq}}\langle \sigma v \rangle$ for the scattering process $\psi H_L \rightarrow N Z_L$ and the Hubble expansion rate $H$ with the temperature of the Universe $T$. In this plot, we have considered the masses of all the components of $H_L$ and $\psi$ to be equal and we have kept them fixed at 1 TeV. From this plot one can easily notice that the ratio $n_{\text{eq}}\langle \sigma v \rangle / H$ is maximum when $T \sim 1$ TeV, same as the mass of incoming particles. After that, $n_{\text{eq}}\langle \sigma v \rangle / H$ reduces with $T$ as the equilibrium number density becomes exponentially suppressed for $T < 1$ TeV. We have plotted the variation of $n_{\text{eq}}\langle \sigma v \rangle / H$ with respect to $T$ for three different values of Yukawa coupling e.g. $Y_\psi = 10^{-4}$, $10^{-6}$ and $10^{-7}$ respectively. We have found that for $Y_\psi \geq 10^{-6}$, our DM candidate remains in thermal equilibrium with the plasma through the scattering $\bar{\psi}(t) H_{L(R)} \rightarrow N Z_{L(R)}$, as the interaction rate exceeds the Hubble expansion rate for the considered range of Temperature $10^2$ GeV $\leq T \leq 10^4$ GeV. However, as we reduce the Yukawa coupling further, the corresponding interaction rate also decreases ($\propto Y_\psi^2$) and we have found that for $Y_\psi \leq 10^{-7}$, $N$ never thermalises with the bath particles. Therefore, the
non-thermality condition \( n_{\text{eq}} \langle \sigma v \rangle / H < 1 \) demands \( Y_\psi, Y_{\psi'} \leq 10^{-7} \).

Nevertheless, one can still consider \( N \) as a non-thermal dark matter candidate without assuming such a low value for the Yukawa couplings. This is in fact, the prime motivation of this article, to realise FIMP dark matter without highly fine-tuned dimensionless couplings. In that case, we have to consider a scenario with low reheat temperature \([58]\) of the Universe so that these particles \( H_{L,R}, \psi, \psi' \) were not thermally produced in the early Universe, which actually prevents \( N \) to thermalise via scattering with \( H_L, H_R, \psi \) and \( \psi' \). For example, the authors of \([59]\) considered such heavy mediators having mass greater than the reheat temperature, but in a different dark matter scenario. Since the masses of these particles appear as bare mass terms in the Lagrangian and hence they do not depend upon the scale of symmetry breaking, we can push them high to any scale above the scale of reheating. Therefore, we do not really have to rely upon very specific inflationary models where low reheat temperature occurs. Such a setup leads to a situation where \( N \) can ‘talk’ to other particles in the thermal bath only through the one loop suppressed portals like \( Z_L, Z_R \) and \( \Phi \). In the right panel of Fig. 3, we plot the ratio of thermally averaged partial decay width of \( Z_{L,R} \) into two \( N \)'s and \( H \) with \( T \). Here, green dashed line represents the variation \( \langle \Gamma \rangle / H \) with \( T \).

\[^2\text{The similar bound on } Y_{\psi'} \text{ is coming from scattering involving } \psi', H_R, N \text{ and } Z_R.\]
for the right handed neutral gauge boson $Z_R$. In this plot, we have considered $M_{Z_R} = 6$ TeV and the corresponding Yukawa coupling $Y_{\psi'} = 10^{-3}$ while the similar plot for the left handed neutral gauge boson $Z_L$ has been drawn for $Y_\psi = 3.38 \times 10^{-4}$ and this is represented by the red solid line. From this figure, it is evident that for such moderately large Yukawa couplings (larger compared to those required for scattering to be out of equilibrium) the decay rates of $Z_L \rightarrow \bar{N}N$, $Z_R \rightarrow \bar{N}N$ always lie below the Hubble expansion rate $H$, thus maintaining the non-thermality criteria of $N$. Although, for $Z_R$, the corresponding decay rate exceeds $H$ when $T < 10$ GeV, however at such low temperature the number density of $Z_R$ with mass 6 TeV becomes exponentially suppressed and hence does have a very little impact on $N$. In fact, $\langle \Gamma \rangle_H$ for both $Z_L$ and $Z_R$ lie well below the expansion rate at $T \sim M_{Z_i}$ ($i = L, R$) where the maximum production of $N$ from the decay of $Z_i$ occurs. Therefore, $N$ remains out of thermal equilibrium for moderately large values of Yukawa couplings ($\lesssim \mathcal{O}(10^{-3})$) and can be a candidate for FIMP accordingly. Note that such Yukawa couplings are much less fine tuned than the ones considered in generic FIMP models and also larger than the electron Yukawa coupling in the Standard Model.

The Boltzmann equation which governs the evaluation of comoving number density (ratio of number density to entropy density) of a FIMP is given by

$$
\frac{dY_N}{dz} = \frac{2M_{Pl}}{1.66M_{sc}^2} \sqrt{g_*(z)} \left[ \sum_{\chi = Z_L, Z_R, \Phi} \langle \Gamma_{\chi \rightarrow \bar{N}N} \rangle (Y_{eq}^\chi - Y_N) \right] + \frac{4\pi^2}{45} \frac{M_{pl}M_{sc}}{1.66} \sqrt{g_*(z)} \left[ \sum_{p = \text{SM fermions}} \left( \sigma v_{p\bar{p} \rightarrow \bar{N}N} \right) \left\{ \left( Y_{eq}^p \right)^2 - Y_N^2 \right\} \right].
$$

(7)

where $z = \frac{M_{sc}}{T}$, is a dimensionless variable while $M_{sc}$ is some arbitrary mass scale which we choose equal to the mass of $Z_L$ and $M_{Pl}$ is the Planck mass. Moreover, $g_*(z)$ is the number of effective degrees of freedom associated to the entropy density of the Universe and the quantity $g_*(z)$ is defined as

$$
\sqrt{g_*(z)} = \frac{g_*(z)}{\sqrt{g_0(z)}} \left( 1 - \frac{1}{3} \frac{d \ln g_*(z)}{d \ln z} \right).
$$

(8)

Here, $g_0(z)$ denotes the effective number of degrees of freedom related to the energy density of the Universe at $z = \frac{M_{sc}}{T}$. The first term in the right hand side of the above Boltzmann equation (7) represents the production of our dark matter candidate $N$ from the decays of
$Z_L$, $Z_R$ and bi-doublet $\Phi$. The quantity $Y_{\chi}^{eq}$ is the equilibrium comoving number density of the species $\chi$ ($\chi = Z_L, Z_R, \Phi$) and in this work, except the dark matter candidate $N$, we consider Maxwell-Boltzmann distribution function for all the other particles which are in thermal equilibrium. As we have mentioned earlier, the production processes of $N$ from the decays of $Z_L$, $Z_R$ and $\Phi$ are one loop suppressed. The Feynman diagrams of these processes are shown in Fig. 1 and the corresponding effective one loop vertices are given in Eq. (6). Using the expressions of effective vertices one can easily compute the thermally averaged decay width for the processes $Z_{L,R} \to \bar{N}N$ and expressions are given by,

$$\langle \Gamma_{Z_{L,R} \to \bar{N}N} \rangle = \frac{K_1 \left( \frac{M_{Z_{L,R}}}{T} \right)}{K_2 \left( \frac{M_{Z_{L,R}}}{T} \right)} ,$$

$$\Gamma_{Z_{L,R} \to \bar{N}N} = \frac{|A_{Z_{L,R}}^{NN}|^2 M_{Z_{L,R}}}{24\pi} \left( 1 - \frac{4 M_N^2}{M_{Z_{L,R}}^2} \right)^{3/2} .$$

Where, $K_n \left( \frac{M_{Z_{L,R}}}{T} \right)$ is the $n$-th order modified Bessel function of second kind. The second term in the right hand side of the Boltzmann equation represents the contributions coming from the annihilations of SM particles to the production processes of $N$. Since $N$ is a singlet under the left-right symmetry group, the interactions of $N$ with the SM particles are possible only through the portal interactions by $Z_L$, $Z_R$ and $\Phi$. One could have a tree level coupling like $\bar{\ell}_L \tilde{H}_L N$ had the additional discrete symmetry $Z_2$ was not in place. The other tree level couplings involving $H_{L,R}, \psi, \psi'$ and $N$ will not play a role in the production of $N$ if these heavy incoming particles were not thermally produced in the early Universe due to their masses being heavier than the reheat temperature after inflation, as argued previously.

Hence, the contributions of such processes are sub-dominant compared to that from the decays of $Z_L$, $Z_R$, $\Phi$ as the couplings between our FIMP dark matter $N$ and these mediator particles are one loop suppressed. In our calculations, we consider only the decay of the neutral gauge bosons $Z_L$, $Z_R$ as they are more likely to be dominant due to the presence of gauge couplings in one of the vertices of the one loop diagram. For $\Phi$ decay diagram, one has more freedom in choosing another Yukawa coupling and hence that contribution can remain suppressed compared to the gauge boson ones.

Finally, the relic density of dark matter, which is defined as the ratio between dark matter mass density and critical density of the Universe, is obtained using the solution of Boltzmann
Production of N from Z\textsubscript{L} decay

Production of N from Z\textsubscript{R} decay

\[ Y_{\psi} = 1.1 \times 10^{-4}, Y_{\psi}' = 1.19 \times 10^{-3} \]
\[ Y_{\psi} = 3.38 \times 10^{-4}, Y_{\psi}' = 1.0 \times 10^{-3} \]
\[ Y_{\psi} = Y_{\psi}' = 4 \times 10^{-4} \]
\[ \Omega_{\text{DM}} h^2 = 0.12 \]

\( \Omega_{\text{DM}} h^2 \) is the relic density of cold dark matter in the early Universe. The figure shows the variation of the relic density with the dimensionless variable \( z = M_{Z_L} / T \) for different values of Yukawa couplings. The plot has been generated for \( M_{Z_R} = 6 \text{ TeV}, M_N = 1 \text{ MeV} \) and three different values of Yukawa couplings \( Y_{\psi} \) and \( Y_{\psi}' \) which reproduce the correct dark matter relic density.

In order to kinematically forbid the tree level decay of either \( H_L (H_R) \) or \( \psi (\psi') \) as well as to use the simple expressions for one loop decay widths in (6) we choose \( M_{H_L} = M_{\psi} = M_L \) (\( M_{H_R} = M_{\psi'} = M_R \)). As mentioned earlier, we have also considered \( M_L, M_R, M, M' \) to be very large at least greater than the reheat temperature of the Universe after inflation, which can be assumed to be as high as \( \sim 10^{10} \text{ GeV} \) in some cosmological scenarios like for example, [61, 62]. As a result, the production of N from the scatterings of these heavy particles is not efficient.

From this plot in Fig. 4 it is seen that at first the relic density of N increases as \( z \) increases from 0.01 to 0.1 (corresponding temperature decreases from \( M_{Z_L} / 0.01 \text{ GeV} \) to \( M_{Z_L} / 0.1 \text{ GeV} \)) and then for \( 0.1 \leq z \leq 0.35 \), \( \Omega_{\text{DM}} h^2 \) saturates to a particular value which depends upon the value of Yukawa coupling \( Y_{\psi'} \) (e.g. green dashed-dotted line and red solid line). This initial rise in the relic density of N (or equivalently in \( Y_N \)) is due to its production from the heavy
right handed neutral gauge boson $Z_R$. Since $Z_R$ is in equilibrium with the thermal bath, most of the production of $N$ from the decay of $Z_R$ occurs for the temperature of the Universe $T \sim M_{Z_R}$. As the temperature drops below the mass of $Z_R$, the number density of $Z_R$ starts becoming exponentially suppressed (Boltzmann suppressed) and finally for $z \geq 0.1$ (or $T \leq 1$ TeV) there are practically not enough number of right handed neutral gauge boson left to produce $N$ and thus the comoving number density and hence the dark matter relic density saturates. Thereafter, again there is a sharp increase in the relic density of $N$ between $z = 1.0$ to $z = 10$ (e.g. red solid line and blue dotted line). This increment of relic density is due to the substantial production of $N$ from the decay of left handed neutral gauge boson $Z_L$ (the usual Z boson in the SM), which depends on the other Yukawa coupling $Y_\psi$ (see Eqs. 6 and 10). Like the previous production regime of $N$ from $Z_R$, in this case also the dominant production of $N$ from $Z_L$ decay occurs at around the temperature $T \sim M_{Z_L}$ ($z \sim 1$). Finally, when the temperature of the Universe drops well below the mass of $Z_L$, all the production modes of $N$ cease and relic density saturates to $\Omega_{DM} h^2 \sim 0.12$, the canonical value measured by the Planck satellite [5]. Here, we have chosen three combinations of Yukawa couplings $Y_\psi$ and $Y_\psi'$ which result in the correct relic density of dark matter. For $Y_\psi = 3.38 \times 10^{-4}$ and $Y_\psi' = 10^{-3}$, we have a situation where there is an equal contribution of both $Z_R$ and $Z_L$ to $\Omega_{DM} h^2$ and that is represented by the red solid line in Fig. 4. On the other hand by tuning both $Y_\psi$ and $Y_\psi'$, one can have scenarios where the production of $N$ is dominated by either $Z_R$ or $Z_L$. These are described by green dashed-dotted line and blue dotted line respectively while the corresponding Yukawa couplings are $Y_\psi = 1.1 \times 10^{-4}$, $Y_\psi' = 1.19 \times 10^{-3}$ and $Y_\psi = Y_\psi' = 4.0 \times 10^{-4}$ respectively.

In Fig. 5, we demonstrate the region in $Y_\psi - Y_\psi'$ plane which is allowed by the observational value of dark matter relic density. Here, we have also varied the mass of $M_{Z_R}$ in the range of 1 TeV to 100 TeV and it has been shown by the colour code. For large value of $Y_\psi'$ i.e $Y_\psi' \gtrsim 10^{-3}$, the production of $N$ is dominated by the $Z_R$ decay and on the other hand, the relic density of $N$ receives maximum contribution from $Z_L$ when $Y_\psi \gtrsim 3.5 \times 10^{-4}$. In the intermediate region, both $Z_L$ and $Z_R$ are contributing to the relic density of $N$ depending upon the values of $Y_\psi$ and $Y_\psi'$. Moreover, we also get a narrow horizontal line for $Y_\psi \sim 4 \times 10^{-4}$ and $Y_\psi' \lesssim 0.7 \times 10^{-3}$, where entire production of $N$ occurs from $Z_L$ decay. The opposite situation where the entire $N$ is produced from $Z_R$ decay occurs for $Y_\psi \sim 1 \times 10^{-4}$ and $Y_\psi' \gtrsim 1.0 \times 10^{-3}$. However, unlike to the $Z_L$ dominated case, here we get a wide band
FIG. 5: Allowed region in $Y_\psi - Y_{\psi'}$ plane which reproduces correct dark matter relic density for $1 \text{ TeV} \leq M_{Z_R} \leq 100 \text{ TeV}$.

and it due to the variation of mass of $Z_R$, which has been varied between 1 TeV to 100 TeV.

Finally, we have also shown the variation of Yukawa couplings $Y_\psi$ and $Y_{\psi'}$ with the mass of our dark matter candidate $N$ in Fig. 6. In this plot, we have varied $M_N$ in the range of 1 keV to 1 TeV. The corresponding variation of $Y_{\psi'}$ ($Y_\psi$) to obtain the correct dark matter relic abundance is indicated by green (dark red) coloured contour. From this plot, one can notice that the allowed values of Yukawa couplings are decreasing with the increase of $M_N$. This can be understood from Eq. 11 which states that for a fixed value of $\Omega_{DM} h^2$ the product of $Y_N(t_0)$ and $M_N$ is a constant. Since, our dark matter candidate is a FIMP, $Y_N(t_0)$ is proportional to the Yukawa couplings. Therefore, to remain within the observed relic density band, any increment in $M_N$ must be accompanied by a decrement in $Y_N(t_0)$ and consequently in $Y_\psi$ and $Y_{\psi'}$. In the right most corner of $M_N - Y_\psi$ plane, one can see that the entire range of $Y_\psi$ is allowed for $M_N > M_{Z_L}/2$. This is due to the fact that in this mass range of $N$, the production from $Z_L$ decay is kinematically forbidden and hence any variation in $Y_\psi$ does not affect the the relic density of $N$. On the other hand, as $Z_L$ decay to a pair of $N$ is not possible for $M_N > M_{Z_L}/2$, the decay $Z_R$ is solely responsible for the entire production of $N$. Therefore, in this mass range of $M_N$, we get a very narrow allowed values
of $Y_{\psi'}$. Moreover, as the entire considered mass range of $N$ is allowed for some combinations of $Y_{\psi}$ and $Y_{\psi'}$, we have checked the nature of our dark matter candidate (hot/warm/cold) by computing its free streaming length following Ref. [63]. We find that $N$ becomes a warm dark matter candidate for $M_N \leq 10$ keV, where its free streaming length goes above 0.01 Mpc. On the other hand, the cold dark matter scenario is viable for $M_N > 10$ keV, where the free streaming length of $N$ always lies below 0.01 Mpc and decreases sharply with the increase of mass of $N$. The possibility of warm dark matter has several motivations, which can be found in the recent review [41].

V. CONCLUSION

We have proposed a UV complete framework to dynamically generate tiny couplings required for non-thermal dark matter scenarios whose relic abundance is generated through the freeze-in mechanism, within the framework called freeze-in massive particle. Based on gauge symmetric extensions of the Standard Model, we particularly consider the left-right symmetric model which have several other motivations related to the origin of parity violation, neutrino mass among others. Considering the dark matter candidate to be a
gauge singlet fermion which has no tree level couplings with the Standard Model particles, we generate its effective couplings with the Standard Model particles at one loop, mediated by gauge bosons and Higgs. After showing such a dark matter candidate to remain out of thermal equilibrium in the early Universe for generic choices Yukawa couplings and masses of particles inside loops, we then calculate its relic abundance by considering both decay and scattering contributions in a way similar to a generic FIMP dark matter candidate. We find that the decay of neutral heavy gauge bosons to a pair of FIMP dark matter candidate is the most dominant production mechanism and can give rise to the correct relic abundance for Yukawa couplings as large as $10^{-4} - 10^{-3}$ while keeping the additional heavy neutral boson mass within experimental reach. Such Yukawa couplings lie in a range which less fine tuned than the electron Yukawa coupling and far less fine tuned than the ones involved in generic FIMP models. On the other hand, a very wide range of dark matter masses is consistent with the relic abundance criteria and some portion of this allowed range can also give rise to warm dark matter scenarios that have several other motivations from small scale structure point of view. Since our UV complete setup has many other particles that lie in the experimentally accessible range, many associated particles can be probed at ongoing experiments, which we leave for future studies. We also note that such a setup can be realised in other gauge extensions of Standard Model as well which may be relatively simpler than the one presented here as a matter of choice.

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Appendix A: Scalar Potential of the Model

The scalar potential for the minimal LRSM is

$$V(\Phi, \Delta_L, \Delta_R) = V_\mu + V_\Phi + V_\Delta + V_{\Phi\Delta} + V_{\Phi\Delta_L\Delta_R},$$

(A1)
where the bilinear terms in Higgs fields are

\[
V_\mu = -\mu_1^2 \text{Tr} [\Phi^\dagger \Phi] - \mu_2^2 \text{Tr} [\Phi^\dagger \Phi^* + \Phi^\dagger \Phi^*] - \mu_3^2 \text{Tr} [\Delta_L^\dagger \Delta_L + \Delta_R^\dagger \Delta_R].
\] (A2)

The self-interaction terms of \( \Phi \) are:

\[
V_\Phi = \lambda_1 \left[ \text{Tr} [\Phi^\dagger \Phi] \right]^2 + \lambda_2 \left[ \text{Tr} [\Phi^\dagger \Phi^*] \right]^2 + \lambda_2 \left[ \text{Tr} [\Phi^\dagger \Phi^*] \right]^2
+ \lambda_3 \text{Tr} [\Phi^\dagger \Phi] \text{Tr}[\Phi^\dagger \Phi^*] + \lambda_4 \text{Tr} [\Phi^\dagger \Phi^*] \text{Tr}[\Phi^\dagger \Phi^*].
\] (A3)

and the \( \Delta_{L,R} \) self- and cross-couplings are as follows:

\[
V_\Delta = \rho_1 \left( \left[ \text{Tr} [\Delta_L^\dagger \Delta_L] \right]^2 + \left[ \text{Tr} [\Delta_R^\dagger \Delta_R] \right]^2 \right) + \rho_3 \text{Tr} [\Delta_L^\dagger \Delta_L] \text{Tr}[\Delta_R^\dagger \Delta_R]
+ \rho_2 \left( \text{Tr} [\Delta_L^\dagger \Delta_L] \text{Tr}[\Delta_L^\dagger \Delta_L] + \text{Tr} [\Delta_R^\dagger \Delta_R] \text{Tr}[\Delta_L^\dagger \Delta_R] \right)
+ \rho_4 \left( \text{Tr} [\Delta_L^\dagger \Delta_L] \text{Tr}[\Delta_R^\dagger \Delta_R] + \text{Tr} [\Delta_L^\dagger \Delta_R] \text{Tr}[\Delta_R^\dagger \Delta_L] \right). \] (A4)

In addition, there are also \( \Phi - \Delta_L \) and \( \Phi - \Delta_R \) interactions present in the model,

\[
V_{\Phi\Delta} = \alpha_1 \text{Tr} [\Phi^\dagger \Phi] \text{Tr}[\Delta_L^\dagger \Delta_L + \Delta_R^\dagger \Delta_R] + \alpha_3 \text{Tr} [\Phi^\dagger \Phi^* \Delta_L^\dagger \Delta_L + \Phi^\dagger \Phi^* \Delta_R^\dagger \Delta_R]
+ \{ \alpha_2 e^{i\beta_2} \text{Tr} [\Phi^\dagger \Phi^*] \text{Tr}[\Delta_L^\dagger \Delta_L] + \alpha_2 e^{i\beta_2} \text{Tr} [\Phi^\dagger \Phi^*] \text{Tr}[\Delta_R^\dagger \Delta_R] \} + \text{H.c.} \) (A5)

with \( \delta_2 = 0 \) making CP conservation explicit, and the \( \Phi - \Delta_L - \Delta_R \) couplings are

\[
V_{\Phi\Delta_{L,R}} = \beta_1 \text{Tr} [\Phi^\dagger \Delta_L^\dagger \Phi \Delta_R + \Delta_R^\dagger \Phi^\dagger \Delta_L \Phi] + \beta_2 \text{Tr} [\Phi^\dagger \Delta_L^\dagger \Phi^* \Delta_R + \Delta_R^\dagger \Phi^* \Delta_L \Phi]
+ \beta_3 \text{Tr} [\Phi^\dagger \Delta_L^\dagger \Phi^* \Delta_R + \Delta_R^\dagger \Phi^* \Delta_L \Phi]. \] (A6)

The scalar potential involving the newly introduced scalar fields beyond the minimal LRSM is

\[
V_{\text{new}} = V_H + V_\sigma + V_{\Phi H} + V_{\Delta H}. \] (A7)

The details of different terms on the right-hand side of the above equation can be written as follows,

\[
V_H = \mu_H^2 \left( H_L^\dagger H_L + H_R^\dagger H_R \right) + \rho_5 \left( \left[ H_L^\dagger H_L \right]^2 + \left[ H_R^\dagger H_R \right]^2 \right)
+ \rho_6 \left[ H_L^\dagger H_L \right] \left[ H_R^\dagger H_R \right], \] (A8)

\[
V_\sigma = \frac{\mu_\sigma^2}{2} \sigma^2 + \rho_8 \sigma^4 + \mu_{\sigma\Delta} \sigma \left( \text{Tr} [\Delta_L^\dagger \Delta_L - \Delta_R^\dagger \Delta_R] \right) + \mu_{\sigma H} \sigma \left( H_L^\dagger H_L - H_R^\dagger H_R \right)
+ \rho_9 \sigma^2 \text{Tr} [\Phi^\dagger \Phi] + \rho_{10} \sigma^2 \left( \text{Tr} [\Delta_R^\dagger \Delta_R + \Delta_L^\dagger \Delta_L] \right) + \rho_{11} \sigma^2 \left( H_L^\dagger H_L + H_R^\dagger H_R \right), \] (A9)
\[ V_{\Phi H} = \mu_{14} H_L^\dagger \Phi H_R + f_{145} \text{Tr}[\Phi^\dagger \Phi](H_L^\dagger H_L + H_R^\dagger H_R), \quad (A10) \]

\[ V_{\Delta H} = (\mu_{15} H_L \Delta_L H_L + \mu_{16} H_R \Delta_R H_R + \text{H.c.}) + f_{145} (\text{Tr}[\Delta_L^\dagger \Delta_L] + \text{Tr}[\Delta_R^\dagger \Delta_R])(H_L^\dagger H_L + H_R^\dagger H_R) \quad (A11) \]

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