



CHARACTERIZING DARK MATTER DYNAMICS AND COLLIDER IMPLICATIONS

THESIS

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To my grandfather and his stories



Declaration



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I hereby declare that the papers included in this thesis are the result of my own work carried out at the Department of Physics, Indian Institute of Technology Guwahati, Guwahati, Assam 781039, India, under the supervision of **Prof. Poullose Poullose**. This thesis has not been submitted to anywhere else for the award of any degree, diploma or associateship.

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It is to certify that the works contained in this thesis by Ms. Sreemanti Chakraborti (Roll No.156121030), a PhD student of the Department of Physics, Indian Institute of Technology Guwahati, Guwahati, Assam 781039, India, are carried out under my supervision and have not been submitted elsewhere for the award of any degree.

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Acronyms

BDT	:	Boosted Decision Tree
BP	:	Benchmark Point
BSM	:	Beyond the Standard Model
CE	:	Chemical Equilibrium
CM	:	Centre-of-Mass
CMB	:	Cosmic Microwave Background
DD	:	Direct Detection
DM	:	Dark Matter
DT	:	Disappearing Tracks
DV	:	Displaced Vertex
FIMP	:	Feebly Interacting Massive Particle
FSR	:	Final State Radiation
HSCP	:	Heavy Stable Charged Particles
ID	:	Indirect Detection
ISR	:	Initial State Radiation
LHC	:	Large Hadron Collider
LLP	:	Long-lived Particle
MCDM	:	Multi-component Dark Matter
MET	:	Missing Transverse Energy
SM	:	Standard Model
SUSY	:	Supersymmetry/Supersymmetric
TMVA	:	Toolkit for Multivariate Analysis
WIMP	:	Weakly Interacting Massive Particle



Abstract

After the Higgs discovery, the Large Hadron Collider (LHC) did not find any evidence of new physics. Despite the Standard Model (SM) being “the” theory for the elementary particles, there are still many unanswered questions that give a strong motivation for beyond the SM (BSM) theories. A particle candidate for dark matter (DM) is one such motivation, as its existence is beyond any doubt from various cosmological and astrophysical observations. In its pursuit, many theories have been proposed for DM origin, the most popular till date being the freeze-out mechanism, featuring Weakly Interacting Massive Particle (WIMP). Although WIMP can explain the observed DM relic density well, it fails to fit the direct detection constraints. This motivates the need for alternatives, featuring the non-WIMP cases.

This thesis focuses on three such non-WIMP scenarios having interesting implications in the DM dynamics and collider probes. In the first part, we discuss two multipartite WIMP scenarios of electroweak scale DM mass. Here, the interplay of stable DM components of different spins and isospins successfully evade the direct detection constraint, which is otherwise not possible in similar single component WIMP models. In the context of collider probe, we show that as a contrast with traditional cut-based MET analysis strategies in the literature, which was initially developed for SUSY analysis, a hard MET cut at low DM mass is not required in the non-SUSY models.

The second part includes two chapters. The first one focuses on reviving the very constrained scalar singlet DM (SDM) model of electroweak scale by extending the dark sector with additional vector-like dark lepton doublet. This induces DM coannihilation and adds additional channels to pair annihilation and all these alleviate the severe direct detection constraints on standard SDM. We also discuss the possible collider probe with multilepton final states and τ -tagging at the LHC using the ROOT based Boosted Decision Tree (BDT) classifier Toolkit for Multivariate Analysis (TMVA). In the second chapter, we add an extra vector-like singlet in the dark sector of the previous chapter. This simple extension helps to do away with ad-hoc addition of new scale of spontaneous symmetry breaking to lift the degeneracy of the lepton doublet components. Furthermore, the mixing between the same-charge dark leptons become a key factor to distinguish between the dark leptons of different isospins, which has not been discussed much in the literature. We discuss how the mixing leaves a positive impact on various search prospects of DM such as indirect detection and collider probe. The proper choice of final states in the LHC feature this mixing effect substantially and a significance study is done for future luminosities.

In the third part, we discuss a minimalistic BSM scenario which identifies different areas of the parameter region for electroweak scale DM mass, where DM relic density is controlled by either non-thermal (freeze-in) or thermal (freeze-out) DM dynamics or their admixture. For suitable parameters, it gives a quantitative estimate of the limiting case where the transition between the above regimes takes place. As a natural consequence of the model arrangement, one or more long-lived particles (LLP) are produced. The DM dynamics often dictates the LLP characteristics and the signatures substantially change from one case to another leading to exotic signatures such as cascade decays of LLPs and kinked charged tracks. The model also features right-handed neutrinos which account for neutrino mass through Type-I seesaw mechanism, that requires it to be long-lived as well. The search prospects for such neutral LLP giving displaced vertex signatures is also discussed for the LHC and future detectors.



List of Publications

- 1) **SREEMANTI CHAKRABORTI** and Rashidul Islam, *Implications of dark sector mixing on leptophilic scalar dark matter*, *J. High Energy Phys.* **03**, 032 (2021).
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- 3) **SREEMANTI CHAKRABORTI**, Victoria Martin and Poulose Poulose, *Freeze-in and freeze-out of dark matter with charged long-lived partners*, *J. Cosmol. Astropart. Phys.* **03**, no.03, 057 (2020).
- 4) **SREEMANTI CHAKRABORTI**, Amit Dutta Banik and Rashidul Islam, *Probing Multicomponent Extension of Inert Doublet Model with a Vector Dark Matter*, *Eur. Phys. J. C* **79**, no.8, 662 (2019).
- 5) **SREEMANTI CHAKRABORTI** and Poulose Poulose, *Interplay of Scalar and Fermionic Components in a Multi-component Dark Matter Scenario*, *Eur. Phys. J. C* **79**, no.5, 420 (2019).



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Introduction

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In this introductory chapter, we intend to set the premise of this thesis: dark matter studies as a window to beyond the Standard Model particle physics. The context of dark matter theory in cosmology and particle physics will be discussed, along with due stress on the phenomenology and the experimental probes.

1.1 Dark Matter: Theory and Phenomenology

The existence of dark matter (DM) in the Universe is one of the most significant open problems in cosmology, particle physics and astrophysics. The first observation of DM dates back to 1933, when Zwicky, in his article [224] studied the motion of the galaxies in the Coma cluster. To account for the discrepancy of mass deficit in his observations, he proposed a non-luminous and non-baryonic component and named it as *dark matter*. Subsequently, there have been several observational pieces of evidence of DM. We enumerate these in the following subsection.

1.1.1 Evidence of Dark Matter

• Rotation Curve of Spiral Galaxies

One of the most convincing pieces of evidence for the existence of DM comes from the phenomena of flatness of rotation curves in the spiral galaxies. Since 1940, there have been observations of anomalies in rotational velocity of the spiral galaxies, mostly in the outer regions. Observed rotation curves usually show that the central part of the galaxy rotates like a rigid body, i.e., $v(r) \propto r$ and then after reaching a maximum value, the velocity saturates.

This is contrary to the classical theories, where it is expected that the velocity towards the outer regions of the galaxy should decrease with the radius. If for simplicity, one considers the spherical distribution of matter in the galaxy, then according to Newtonian dynamics, one should expect that the rotational velocity of an object on a stable orbit with radius r from the centre of a galaxy should be,

$$v(r) \propto \sqrt{\frac{M(r)}{r}}$$

where $M(r) = 4\pi \int_0^r \rho(a) a^2$ is the mass of matter having density $\rho(a)$ inside the sphere of radius r . So, outside the visible disk, M should be constant and $v \propto 1/\sqrt{r}$.

Contrary to this, the actual observation of the flat behaviour of the velocity, i.e., $v(r) \simeq \text{constant}$ suggests the existence of a *dark halo* outside the visible disk with $M(r) \propto r$ and $\rho(r) \propto 1/r^2$. Also, this further indicates that a large part of the mass is present in the outer part of the galaxy beyond the visible disk. Such measurement of the rotation curves was first made by Ford and Rubin in 1978 [208].

• **Bullet Clusters**

Bullet cluster is formed after the collision of two clusters of galaxies. The concentration of mass in the cluster is determined using the effect of gravitational lensing, where light from distant objects bends around the gravitating matter in the cluster. It is found that most of the ordinary visible matter in the cluster is separate from the matter responsible for observed gravitational lensing, signifying that nearly all of the matter in the cluster is non-visible, i.e., dark.

• **Anisotropies in CMB**

The amount of DM in the Universe can be extracted from the analysis of the Cosmic Microwave Background Radiation (CMB). CMB is known to be isotropic with temperature variation closely similar to black body radiation of temperature $T = 2.728$ K, but it consists of tiny temperature fluctuations called anisotropy. The variations of temperature of CMB can be expressed as the sum of spherical harmonic Y_{lm} , with the power spectrum against the angular momentum representing various information about the constituent of the Universe. Experiments like WMAP [183], PLANCK [16] and COBE [52] measure the height of the Doppler peaks in CMB power spectrum and the latest PLANCK data estimates that the Universe consists of only around 4.9% of visible baryonic matter while about 26.8% is non-baryonic dark matter and the rest 68.3% is dark energy.

1.1.2 Dark Matter from Particle Physics Perspective

Although the exact nature of DM is not known till now, there have been many studies proposing DM candidates over a huge mass range from 10^{-22} eV (fuzzy dark matter) to 10^{48} GeV (primordial black hole). DM as an elementary particle is also one of the most popular hypotheses explored so far. To fit the bill as an ideal particle DM candidate, the following properties as known from the observations need to be fulfilled [182]

- *Stability* : DM should not decay to any visible particle, so it should be stable over the age of the Universe.
- *Charge neutrality* : DM should be dark, i.e., should not interact with photons. So it must be charge neutral.
- *Weakly interacting* : Apart from gravitational interaction, DM only has a very weak interaction with the visible matter.
- *Massive* : DM particles must be massive so that they contribute substantially to the energy density of the Universe, leading to structure formation.
- *Non-baryonic* : The baryonic matter which is massive and small is mostly gas and dust. Both of these things, especially when hot, are easily observable. Therefore, DM is traditionally hypothesized as non-baryonic and having no strong interaction.

Based on the above properties, a broad classification of DM candidates exist

- **Cold dark matter** : Cold dark matter (CDM) signifies non-relativistic DM particles. Due to their relatively slow velocity at the matter-radiation decoupling time in the early Universe, it is assumed that CDM formed galactic structure in a bottom-up approach which agrees with the experimental evidence. Candidates for CDM include WIMPs (Weakly Interacting Massive Particle), axions, MACHOs (Massive Astrophysical Compact Halo Object) (although non-baryonic) etc.

• **Hot dark matter** : Hot dark matter (HDM) consists of relativistic particles. The primary candidate for HDM is the neutrino. However, HDM theory does not account for the galactic structure formation observed in the Universe. Therefore, any pure HDM model of the Universe is ruled out. In the literature, HDM is discussed in the context of Mixed Dark Matter models, where the bulk of the dark matter is cold, but a tiny fraction is hot.

WIMPs and beyond : WIMPs are an example of non-baryonic CDM candidate. They interact with visible matter only via the weak interaction. For a stable WIMP of electroweak scale, i.e., GeV-TeV scale mass, the weak scale cross-section ($\langle\sigma v\rangle \sim 10^{-26}\text{cm}^3\text{s}^{-1}$) of DM-visible matter interaction produces the right relic density consistent with the observations. This is known as the WIMP miracle. WIMP miracle provides a model independent interpretation of DM at the weak scale. WIMP includes neutralino [129, 147], gravitino [80], which appear in the weak scale supersymmetric (SUSY) models. Examples of a few non-SUSY WIMP candidates are scalar singlet DM [196], inert doublet CP even scalar [188], fermionic singlet DM [180], vector boson DM (VBDM) [133].

WIMPs are one of the most popular DM candidates in the electroweak scale DM phenomenology [207]. However, WIMPs are severely constrained by the latest experimental searches. This leads to proposing non-WIMP DM candidates. We will elaborate this in the following subsections.

1.1.3 Possible Theories of Dark Matter Evolution

Thermal Freeze-Out of Single-Component WIMP

Dark matter candidates are categorized according to whether they are produced as “thermal” or “non-thermal” relics. According to the theory of thermal relics, the early Universe was initially dense and hot, and all the particles were in thermal equilibrium. Afterwards, along with expansion, the Universe began to cool down. When the rate of interaction among DM and other particles in the thermal bath could not compete with the rate of expansion of the Universe, then the DM particles decoupled from the thermal bath and the number of DM particles per co-moving volume asymptotically approached a constant. This is known as “thermal freeze-out” and the co-moving number density came to be known as relic density, or relic abundance. The standard calculation for the thermal freeze-out of WIMPs is given below [145].

The evolution of the phase space distribution function, $f(p, x)$ is governed by the Boltzmann equation

$$L[f] = C[f] \quad (1.1)$$

where L denotes the Liouville operator and C is the collision operator, representing the interactions of the particle species considered.

Quantitatively, the evolution of thermal DM particle of number density n interacting with the particles in a thermal bath is given by the Boltzmann equation

$$\frac{dn}{dt} + 3Hn = -\langle\sigma v\rangle (n^2 - n_{eq}^2) \quad (1.2)$$

where t is time, H is the Hubble constant, ($H \equiv \dot{a}/a$, where a is the scale factor of the Universe). σ is the cross-section for DM pair annihilation, v is the relative speed of annihilating DM particle and $\langle\sigma v\rangle$ denotes the thermal averaged annihilation cross section, signifying an average over the thermal distribution of momenta of DM particles. Also, the spin average should be average over the initial states and sum over the final states. n_{eq} denotes equilibrium number density. In the non-relativistic limit and the Maxwell-Boltzmann approximation, it is given by the relation

$$n_{eq} = g \left(\frac{mT}{2\pi} \right)^{3/2} e^{-m/T} \quad (1.3)$$

where g is the degree of freedom of DM species, m is DM mass and T is the temperature of the thermal bath. Furthermore, it is assumed for simplicity that DM can only annihilate and be created in pairs and that there is no particle-antiparticle asymmetry, i.e., $n_{DM} = n_{\overline{DM}} = n$.

Now, the thermal relic density (Ω) is determined by solving Boltzmann equation numerically. For this, we introduce the variables

$$Y \equiv \frac{n}{s}, \quad Y_{eq} \equiv \frac{n_{eq}}{s}$$

where s denotes the entropy density given by $s = 2\pi^2 g_* T^3/45$, g_* is given by the number of relativistic degrees of freedom. From the conservation of entropy per co-moving volume, one obtains

$$\dot{n} + 3Hn = s\dot{Y}, \quad s\dot{Y} = -\langle\sigma v\rangle s^2 (Y^2 - Y_{eq}^2)$$

Also, we obtain the following expression on further substituting $x \equiv m/T$

$$\frac{dY}{dx} = -\frac{\langle\sigma v\rangle s}{Hx} (Y^2 - Y_{eq}^2) \quad (1.4)$$

$\langle\sigma v\rangle$ can also be written in the form of non-relativistic expansion in powers of v^2

$$\sigma v = a + bv^2 + \mathcal{O}(v^4) \simeq a + 6\frac{b}{x} \quad (1.5)$$

where a, b are constants. For s -wave annihilation, a gives the dominant contribution to the annihilation of WIMPS. If s -wave contribution is suppressed, b is described by the p -wave contribution.

The above approximation leads to the following relation in terms of the variable $\Delta = Y - Y_{eq}$

$$\Delta' = -Y_{eq}' - f(x) \Delta (2Y_{eq} + \Delta) \quad (1.6)$$

where prime denotes $\frac{d}{dx}$ and $f(x) = \sqrt{\pi g_*/45} m M_{Pl} (a + 6b/x)x^{-2}$. In order to analytically solve Eq. (1.6), the variable $x_F = m/T_F$ is introduced, where T_F is known as the freeze-out temperature. Therefore, we obtain the solution of Eq. (1.6) for two extreme cases $x \ll x_F$ and $x \gg x_F$ as

$$\begin{aligned} \Delta &= -\frac{Y_{eq}'}{2f(x) Y_{eq}} && \text{for } x \ll x_F \\ \Delta' &= -f(x) \Delta^2 && \text{for } x \gg x_F \end{aligned}$$

Integrating the above equation between x_F and ∞ and assuming $\Delta_{x_F} \gg \Delta_\infty$, we obtain the following relation

$$Y_\infty^{-1} = \sqrt{\frac{\pi g_*}{45}} m M_{Pl} x_F^{-1} (a + 3b/x_F) \quad (1.7)$$

where M_{Pl} is the (non-reduced) Planck mass and the expression for the relic density is given by

$$\Omega h^2 = \rho_X / \rho_c \quad (1.8)$$

where $\rho_X = m_X n = m_X s_0 Y_\infty$, s_0 is denoted by the present entropy density of the Universe, given by $s_0 = 2\pi^2 g_* T_0^3/45$. ρ_c is the critical energy density of the present Universe, given by $\rho_c = 3H_0^2/8\pi G$. Substituting the values, one can obtain

$$\Omega h^2 \approx \frac{1.07 \times 10^9 \text{ GeV}^{-1}}{M_{Pl}} \frac{x_F}{\sqrt{g_*}} \frac{1}{(a + 3b/x_F)}$$

where h is the reduced Hubble parameter ($= H_0/100 \text{ km s}^{-1} \text{ Mpc}^{-1}$) and a, b are expressed in GeV^{-2} . x_F can be calculated by the iterative solution of the equation

$$x_F = \ln \left[c(c+2) \sqrt{\frac{45}{8}} \frac{g}{2\pi^3} \frac{m M_{Pl} (a + 6b/x_F)}{g_*^{1/2} x_F^{1/2}} \right] \quad (1.9)$$

where c is a constant of order one and is determined by matching the late-time and early-time solutions. From Eq. (1.9), the relic density can also be approximately evaluated as

$$\Omega h^2 \approx \frac{3 \times 10^{-27} \text{ cm}^3 \text{ s}^{-1}}{\langle \sigma v \rangle} \quad (1.10)$$

In any BSM model accounting for the observed DM relic density, usually, the stability of DM is ensured with the help of symmetry transformation, in which the visible sector particles are invariant, whereas the dark particles undergo a non-trivial transformation. In many models, there is more than one particle belonging to this sector (the dark sector), with the lightest one as the DM particle [207].

Coannihilation: An Exception in Thermal Freeze-out

In coannihilation, one or more dark sector particles have a very small mass split between themselves and the DM. In that case, the standard relic calculation fails [148].

For a generic coannihilation framework, let us consider N particles $X_i (i = 1, \dots, N)$ with masses m_i and internal degrees of freedom g_i . Also, it is assumed that $m_1 \leq m_2 \leq \dots \leq m_{N-1} \leq m_N$ and the lightest particle is protected against decay due to some underlying symmetry. If the lightest particle is denoted by X_i , then the Boltzmann equation becomes

$$\frac{dn}{dt} = -3Hn - \sum_{i,j=1}^N \langle \sigma_{ij} v_{ij} \rangle (n_i n_j - n_i^{eq} n_j^{eq}) \quad (1.11)$$

where n is the number density of DM and $n = \sum_{i=1}^N n_i$. This is because the decay rate of the dark sector particles except DM is much faster than the age of the Universe. The coannihilation cross-section is given by

$$\sigma_{ij} = \sum_X \sigma(X_i X_j \rightarrow X_{SM}) \quad (1.12)$$

Also, the expression for the relative particle velocity is given by

$$v_{ij} = \frac{\sqrt{(p_i \cdot p_j)^2 - m_i^2 m_j^2}}{E_i E_j} \quad (1.13)$$

where p_i and E_i are the four-momentum and energy of the particle i respectively. The thermal average $\langle \sigma_{ij} v_{ij} \rangle$ is defined as

$$\langle \sigma_{ij} v_{ij} \rangle = \frac{\int d^3 p_i d^3 p_j f_i f_j \sigma_{ij} v_{ij}}{\int d^3 p_i d^3 p_j f_i f_j} \quad (1.14)$$

where f_i are distribution functions in the MB distribution. Since the scattering rate of the dark sector particles off particles the thermal bath is much faster than their annihilation rate, we obtain

$$\frac{dn}{dt} = -3Hn - \langle \sigma_{\text{eff}} v \rangle (n^2 - n_{eq}^2) \quad (1.15)$$

where $\langle \sigma_{\text{eff}} v \rangle = \sum_{i,j} \langle \sigma_{ij} v_{ij} \rangle \frac{n_i^{eq}}{n^{eq}} \frac{n_j^{eq}}{n^{eq}}$. It is further convenient to write

$$\frac{n_i^{eq}}{n^{eq}} = \frac{g_i (1 + \Delta_i)^{3/2} e^{-x \Delta_i}}{g_{\text{eff}}}$$

where $\Delta_i = (m_i - m_{DM})/m_{DM}$ and $g_{\text{eff}} = \sum_{i=1}^N g_i (1 + \Delta_i)^{3/2} e^{-x \Delta_i}$.

To summarize the above, in coannihilation, the efficient interaction rate within the dark sector maintains the chemical equilibrium between the DM and heavier dark particles. The dark sector decouples from the SM much before the decoupling within the dark sector occurs. The heavier dark particles decay into DM and when this decay rate falls below the expansion rate of the Universe, DM abundance saturates.

Other Exceptions in Thermal Freeze-out

• **Annihilation near a pole :** The annihilation cross-section of a DM particle X through exchange of the intermediate particle of mass M_{ex} is given by [148, 195]

$$\sigma v = \frac{\alpha_D \alpha_V s}{(s - M_{ex}^2)^2 + M_{ex}^2 \Gamma_{ex}^2}$$

where α_D, α_V corresponds to the coupling of the exchange particle to the dark sector (D) and the visible sector (V) respectively. Γ_{ex} is the total width of the exchange particle and $s = \frac{4m_X^2}{1-v^2/4}$.

On further defining the following quantities,

$$r = \left(\frac{2m_X}{M_{ex}} \right)^2, \quad \epsilon = \left(\frac{\Gamma_{ex}}{M_{ex}} \right)^2$$

one obtains

$$\sigma v = \frac{\alpha_D \alpha_V}{M_{ex}^2} \frac{\frac{r}{1-v^2/4}}{\left[1 - \frac{r}{1-v^2/4} \right]^2 + \epsilon} \quad (1.16)$$

To calculate the relic abundance, the annihilation cross-section needs to be thermally averaged. The standard prescription is to Taylor expand to the first order in v^2 . But, for small ϵ in Eq. (1.16), the expansion in v^2 breaks down near the pole and the standard freeze-out calculation fails.

There are different ways to address this problem. The standard method is to numerically calculate the thermal averaged cross-section from the mean of σv in the CM frame ($\vec{p}_1 = -\vec{p}_2$)

Taking the nonrelativistic approximation, one can obtain

$$\langle \sigma v \rangle = \langle \sigma v \rangle_{\text{num}} = \frac{\int p^2 \langle \sigma v \rangle e^{-p^2/mT} dp}{\int p^2 e^{-p^2/mT} dp} = \frac{x^{3/2}}{2\sqrt{\pi}} \int_0^\infty v^2 \langle \sigma v \rangle e^{-xv^2/4} dv$$

• **Forbidden Freeze-out :** This exception holds when $2m_X \lesssim m_1 + m_2$, where m_X is the DM particle and m_1, m_2 are two SM particles. The annihilation channel $XX \rightarrow 12$ is forbidden in zero temperature, but in the early Universe when the temperature at the decoupling time $T \simeq m_X/25$ is sufficient for s to reach the threshold $m_1 + m_2$, then these channels are allowed. If the annihilation products are not too heavy, then these “forbidden” annihilations can dominate the total DM annihilation cross-section [106, 148, 195].

Thermal Freeze-Out of Multi-Component WIMP

In multi-component DM (MCDM) models [87, 223], there are more than one stable dark sectors with the possibility of interaction between each other. The number density of DM is determined by solving the coupled Boltzmann equation in such cases. As a typical example of a minimalistic multi-component set-up, let us consider two stable DM candidates X_1 and X_2 ($m_{X_2} \geq m_{X_1}$) belonging to dark sectors 1 and 2 respectively. Ideally, one should consider all types of number changing processes such as

- i) *Annihilation* : $X_1 X_1 \rightarrow SM SM, X_2 X_2 \rightarrow SM SM$
- ii) *Semi-annihilation* : $X_1 X_1 \rightarrow X_1 SM, X_2 X_2 \rightarrow X_2 SM$
- iii) *Conversion* : $X_2 X_2 \rightarrow X_1 X_1$

On assuming the chemical equilibrium between X_1, X_2 , the coupled Boltzmann equation becomes

$$\frac{dn_{X_1}}{dt} + 3Hn_{X_1} = -\langle\sigma v\rangle_{X_1 X_1 \rightarrow SM SM} (n_{X_1}^2 - n_{X_1 eq}^2) + \langle\sigma v\rangle_{X_2 X_2 \rightarrow X_1 X_1} \left(n_{X_2}^2 - \frac{n_{X_2 eq}^2}{n_{X_1 eq}^2} n_{X_1}^2 \right); \quad (1.17)$$

$$\frac{dn_{X_2}}{dt} + 3Hn_{X_2} = -\langle\sigma v\rangle_{X_2 X_2 \rightarrow SM SM} (n_{X_2}^2 - n_{X_2 eq}^2) - \langle\sigma v\rangle_{X_2 X_2 \rightarrow X_1 X_1} \left(n_{X_2}^2 - \frac{n_{X_2 eq}^2}{n_{X_1 eq}^2} n_{X_1}^2 \right). \quad (1.18)$$

where $n_i, i = X_1, X_2$, denotes the number density of DM particles and their equilibrium number densities are expressed as n_{ieq} respectively. Solving for the Eqs. (1.17) and (1.18), one can obtain the relic density contributions from each of the DM candidates of the form

$$\Omega_i h^2 = 2.755 \times 10^8 \frac{m_i}{GeV} Y_i(T_0), \quad i = X_1, X_2, \quad (1.19)$$

where $Y_i = n_i/s$, is the yield of DM candidate obtained at present temperature of Universe T_0 and s is the entropy density of the Universe. The total DM relic density is then obtained by adding individual relic density of both the candidates, i.e.,

$$\Omega_{DM} h^2 = \Omega_{X_1} h^2 + \Omega_{X_2} h^2, \quad (1.20)$$

Therefore, as a brief interpretation of the above expressions, an interacting MCDM framework has an efficient DM-DM conversion, which maintains the chemical equilibrium between the DM components. The individual components decouple from the SM much before the dark sector decoupling occurs. The heavier DM particle undergoes pair conversion into a pair of lighter DM and when this conversion rate falls below the expansion rate of the Universe, the total DM abundance saturates.

Non-Thermal Dark Matter

- **Freeze-in** : Non-thermal DM [54, 154] has a characteristically different thermal history than WIMPs. It is assumed that in the early Universe, DM has a negligible abundance and has a very feeble interaction with the particles in the thermal bath so that it is decoupled from the thermal plasma. DM number density changes as DM is gradually produced from the decay or scattering of the particles in the thermal bath. When the temperature of the Universe drops below the DM mass (if DM is heavier than the bath particles) or the mass of the decaying thermal bath particle (if DM is lighter than the decaying bath particle), the DM abundance saturates, or “freezes-in” with a yield that increases with the interaction strength of DM with the thermal bath. Typically, the interaction strength required to get the right relic density by freeze-in is orders of magnitude smaller than that of WIMP.

Non-thermal production can occur from both decay and scattering of the bath particles. For simplicity of calculation, we only discuss the decay production in the next chapters of this thesis. The standard calculation, in that case, is given below [154]

Let us consider a Feebly Interacting Massive Particle (FIMP) X with a negligible coupling with the bath particles A and B in the early Universe. $m_A > m_B + m_X$, and the freeze-in mechanism predominantly occurs through the decay $A \rightarrow B + X$.

The Boltzmann equation can be written as

$$\begin{aligned} \dot{n}_X + 3Hn_X = & \int d\Pi_X d\Pi_A d\Pi_B (2\pi)^4 \delta^4(p_X + p_B - p_A) \\ & \times [|M|_{A \rightarrow B+X}^2 f_A (1 \pm f_B)(1 \pm f_X) - |M|_{B+X \rightarrow A}^2 f_B f_X (1 \pm f_A)] \end{aligned} \quad (1.21)$$

where $d\Pi = d^3p_i/(2\pi)^3 2E_i$ are phase space elements, f_i is the phase space density of the i th particle.

It is assumed that the initial X abundance is negligible. Therefore, $f_X = 0$ and the second term of Eq. (1.21) can be neglected.

Approximating $(1 \pm f_B) \approx 1$ and defining the number density of the i th particle, $n_i = \frac{g_i}{2\pi^3} \int d^3p f_i$, Eq. (1.21) can be rewritten as

$$\dot{n}_X + 3Hn_X \approx 2g_A \int d\Pi_A \Gamma_A m_A f_A \approx g_A \int \frac{d^3p_A f_A \Gamma_A}{(2\pi)^3 \gamma_A} \quad (1.22)$$

where $\gamma_A = E_A/m_A$ and Γ_A is the partial width of $A \rightarrow B + X$ decay.

Assuming kinetic equilibrium for the bath particles, f_A can be approximated as $f_A = (e^{E_A/T} \pm 1)^{-1} \approx e^{-E_A/T}$.

Conversion of integral over momentum space into an integral over energy, the Eq. becomes

$$\begin{aligned} \dot{n}_X + 3Hn_X & \approx g_A \int_{m_A}^{\infty} \frac{m_A \Gamma_A}{2\pi^2} (E_A^2 - m_A^2)^{1/2} e^{-E_A/T} dE_A \\ & = \frac{g_A m_A^2 \Gamma_A}{2\pi^2} T K_1(m_A/T) \end{aligned} \quad (1.23)$$

where K_1 is the modified Bessel function.

Using approximations $\dot{T} \equiv -HT$, one can solve for the yield $Y \equiv n/S$

$$Y_X \approx \int_{T_{min}}^{T_{max}} \frac{g_A m_A^2 \Gamma_A}{2\pi^2} \frac{K_1(m_A/T)}{SH} dT \quad (1.24)$$

where $S = 2\pi^2 g_*^S T^3/45$ and $H = 1.66 \sqrt{g_*^\rho} T^2/M_{Pl}$. $g_*^{S,\rho}$ are the effective degrees of freedom of the bath particles at the freeze-in scale, i.e., $T \sim m_A$ for the entropy and the energy density respectively.

Substituting $x \equiv m/T$ and appropriately considering the reheating temperature, the final yield becomes

$$Y_X \approx \frac{135 g_A}{8\pi^3 (1.66) g_*^S \sqrt{g_*^\rho}} \left(\frac{M_{Pl} \Gamma_A}{m_A^2} \right) \quad (1.25)$$

and the relic density expression becomes

$$\Omega_X h^2 \approx \frac{1.09 \times 10^{27} g_A}{g_*^S \sqrt{g_*^\rho}} \frac{m_X \Gamma_A}{m_A^2} \quad (1.26)$$

In the above calculation, it is assumed that the decaying particle is always in equilibrium with thermal plasma. In principle, it can also freeze-out at some point of time in the history of the Universe and then undergo a late decay to produce the FIMP DM. These are discussed in [154]. However, this is not relevant to our study which we have discussed in details in this thesis later on.

• **Other out-of-equilibrium phenomena :** Assumption of chemical equilibrium (CE) in the context of the dark sector particles can change the evolution of DM number density substantially. This critically depends on the interaction between DM and SM as well as within the dark sector or among more than one dark sectors. As an easy example of this, [172] shows a schematic picture on how the dominant

DM number changing processes vary on changing these interaction strengths from very small (super-WIMP regime) to large enough value within perturbative limits (thermal freeze-out regime). When the interaction within the dark sector drops below the limit where CE is no longer possible to maintain, then one needs to solve the full Boltzmann Eq. taking into account all the possible conversion processes, i.e., forward as well as inverse decay and scattering processes. These out-of-equilibrium phenomena, such as co-scattering [115], mediator-driven annihilation [66, 172] etc. often lead to additional feasible regions in DM phenomenology compared to WIMP or FIMP limit and can be an interesting handle for collider probe of the dark mediators [138, 172, 179] as well. Also, since for such phases, the interaction strength is not so feeble as in the FIMP regime, the collider probe (e.g.: exotic long-lived particle signatures [179]) becomes more interesting than FIMP probe.

1.1.4 Detection Prospects of Dark Matter

The strategies for the detection of dark matter can be as follows-

Direct Detection

Direct detection experiments attempt to observe DM particles as they pass through the Earth. Even though DM is expected to have a very scarce interaction, its density in the galactic halo of the milky way should be sufficiently large so that scattering processes with nuclei occur occasionally.

WIMPs can scatter with ordinary matter through $DM\ SM \rightarrow DM\ SM$ interactions. These scatterings can be elastic, inelastic, spin-independent or spin-dependent. [53, 55, 136]

In the elastic scattering experiments, WIMP scattering off the nucleus of the target material is measured in the form of the recoil energy spectrum in the low background, sensitive detectors. These experiments measure the event rate, i.e., the number of events per unit time per unit mass of detector material as a function of energy deposited in the detector.

In inelastic scattering [55], contrary to the above case, the nuclear recoil is not measured in the target. Several possibilities are considered in this context, such as ionization of the target when WIMP interacts with the bound target electrons. Alternatively, WIMP can excite the target nuclei as well. As an observable signature of these phenomena, the recoil is followed by an immediate (nanosecond later) photon emission.

In this thesis, we will discuss the results from the elastic scattering experiments in the context of our studies. In such experiments, the differential event rate for a target consisting of N_i nuclei and corresponding to the recoil energy E_R is given by [225]

$$\frac{dR}{dE_R} = N_i \frac{\rho_{DM}}{m_{DM}} \int_{v > v_{min}} d^3v f(\vec{v} + \vec{v}_E(t)) \frac{d\sigma_{DM}^i(v, E_R)}{dE_R} \quad (1.27)$$

where $v = |\vec{v}|$ is the speed of the particle in the lab frame, $f(\vec{v} + \vec{v}_E(t))$ is the velocity dispersion in the Earth's frame, v_{min} is the minimum speed of DM particles and σ_{DM}^i is the scattering cross-section of DM particles off the nuclei.

The recoil energy is given by

$$E_R = \frac{m_r^2 v^2 (1 - \cos \theta)}{m_i}$$

where θ is the scattering angle in the CM frame and m_r is the reduced mass of the nucleus-DM system.

The spin dependence in the expression for the differential cross-section comes through the DM-nuclei coupling. Axial-vector interaction [55, 116, 186] is an example of spin-dependent interaction, where the scattering cross-section is proportional to $J(J+1)$ due to coupling to the spin-content of a nucleon. On the other hand, in spin-independent interaction, the cross-section increases with the mass of the target nuclei. Therefore, for heavy atoms, the spin-independent scattering dominates over the spin-dependent counterpart. Scalar interaction is an example of spin-independent scattering.

On assuming the standard spin-dependent and spin-independent interactions, one obtains the time-integrated differential cross-section as

$$\frac{d\sigma_{DM}^i(v, E_R)}{dE_R} = \frac{m_i}{2m_r^2 v^2} (\sigma_0^{SI} F_{SI}^2(E_R) + \sigma_0^{SD} F_{SD}^2(E_R))$$

where $F^2(E_R)$ are nuclear form factors and σ_0 is the cross-sections in the limit of zero momentum transfer.

The total event rate integrated over the recoil energy above a detector dependent threshold can be compared with the upper limits from different direct search experiments. In this way, recent direct detection experiments such as LUX [209], PANDA-II [112], XENON-1T [29] etc. constrain the upper limit of DM-nucleon scattering cross-section as a function of WIMP mass around the electroweak scale, i.e., from a few GeV up to the TeV scale.

As an interesting note, MCDM models leave interesting signatures in the direct detection spectra. MCDM models with large mass splitting between the DM components can give a kink in the recoil spectra for DM-nuclei elastic scattering [61, 122, 160]. However, we will not explore much details of this feature in this thesis.

Indirect detection

Indirect detection experiments attempt to observe the products of $DM DM \rightarrow SM SM$ annihilation which take place in regions of high matter density in the Universe. The annihilation products travel to the Earth without getting much affected and include high energy γ -rays, neutrinos, positrons, muons etc. which have a characteristic energy spectrum. These can be observed with many ground-based and space-based telescopes [53, 55, 136].

The flux of DM annihilation products is proportional to the number of annihilations per unit volume, per unit time, i.e., $\propto \sigma v n^2(r) \equiv \sigma v \rho^2(r)/m_{DM}^2$, where m_{DM} is DM mass, σv is the annihilation cross section multiplied by the velocity, $n(r)$ and $\rho(r)$ are the DM number density and mass density respectively and r is the distance from the galactic centre.

Taking into account all the factors, the observed flux detected by the experiments is given by

$$\Phi_i(\psi, E) = \sigma v \frac{dN_i}{dE} \frac{1}{4\pi m_{DM}^2} \int_{\text{line of sight}} ds \rho^2(r(s, \psi)) \quad (1.28)$$

where the index i denotes the secondary particles observed and the coordinate s runs along the line of sight, in a direction making an angle, ψ , from the galactic centre. $\frac{dN_i}{dE}$ denotes the rate of annihilation of secondary particles, denoted by i .

For WIMP mass around the electroweak scale, the final states for annihilation can include $b\bar{b}$, $\mu^+\mu^-$, $\tau^+\tau^-$, W^+W^- etc. One can constrain the thermally averaged annihilation cross-section for annihilation into different final states by experimentally measuring the flux. Experiments such as HESS [13], FERMI-LAT [20] etc. puts the bound on the upper limit of $\langle\sigma v\rangle$ as a function of WIMP mass.

Collider Signatures

Particle colliders can be a probe for DM as well. DM is produced at the colliders through $SM SM \rightarrow \{SM\} + E_T^{miss}$, where $E_T^{miss} \equiv DM DM + \nu$ and $\{SM\}$ denotes one or more visible SM particles. It is assumed that DM escapes the detector without interacting and leaves the missing energy signature along with the visible signal from the SM final states. These visible signals manifest in the collider as leptonic or jet final states as well as mono- X signatures [216], where X can be a photon, gluon, W or Z coming out of initial state radiation (ISR).

In the case where DM is the only particle in the dark sector of the particle spectrum, the most prominent signature is the mono- X . On the other hand, if the dark sector has additional particles, which can be produced through gauge interaction (or through interaction with Higgs boson), their decay

can give rise to interesting signals beyond mono-X. Depending on the decay of the dark sector particles leading to DM and light SM particles in the final state, two characteristically different collider signatures are obtained

- **Prompt Decay Signatures:** A few typical examples of the prompt decay signatures in DM models can be [225]

1. The transverse momentum imbalance in an event due to the presence of DM particles produced together with one or more SM particle,
2. a bump in the dijet or dilepton invariant mass distributions,
3. an excess in the dijet angular distribution, produced by a DM mediator like Higgs or Z mediator.

There have been numerous searches at the LHC for probing DM signals as a smoking gun for BSM physics. One can refer to these review particles for an overview [24, 82, 135, 173, 198].

- **Delayed Decay Signatures:** Unlike the prompt decay signature where usually WIMPs can be probed, the delayed decay signatures are very useful to probe non-thermal DM [49], compressed DM scenarios [73] etc. For these signatures, the decaying dark sector particle has a much larger decay length than the SM particles, so that they decay within the detector volume at a considerable distance away from the interaction point or even escape the detector. These class of particles are therefore known as the long-lived particles [185]. Although there can be many possibilities of the entire class of these exotic signatures, the recent searches in the LHC focus on the following signatures [45, 139, 185]

1. **Heavy Stable Charged Particles (HSCP):** When the decay length of a charged metastable particle becomes larger than the dimension of the detector, the particle escapes without decaying in the detector, but forms a track in the Inner Detector (ID). If the particle is reasonably heavier than a proton, its speed β will be significantly lower than that of any track-forming SM particle of the same momentum. One way to detect this is via specific ionisation. The average ionisation energy loss per unit distance travelled by a charged particle in the material of specific density, $\langle \frac{dE}{dx} \rangle$, has a functional dependence on β given by the Bethe-Bloch formula [225]. This shows that slower the charged particle, larger is the specific ionisation. Silicon-based and gaseous tracking detectors are typically used to measure this charge deposition associated with a hit.
2. **Disappearing Tracks (DT):** These signatures arise when a charged LLP lives long enough to travel deep into the ID but decays at some point within it into a visible and another invisible particle. To be precise, the track of the incoming LLP seems to disappear midway through the ID when the charged decay product is too soft to be reconstructed. Since DT are necessarily short and have few hits, their momentum-measurement resolution is poorer than that of the HSCP tracks.
3. **Displaced Vertex (DV):** These signatures arise when a neutral LLP undergoes a long decay into a pair of leptons or jets. Since the DM in the final state always gives missing energy, the momenta of the visible displaced decay products do not point to the interaction point.

Since no signal for DM has been yet observed in the LHC so far, limits are set on masses, couplings, and cross-sections.

Collider phenomenology of DM models is essential to identify and characterize the properties of the DM candidates. Even if direct detection measurements discover dark matter particles, their ability to probe their properties is quite limited. On the other hand, colliders like LHC are capable of probing the detailed properties at high precision. This would also lead to model discriminations once the discovery is made.

1.2 Thesis objective

With all the necessary introduction to phenomenology and search prospect of particle DM of electroweak scale, in this thesis, we explore various non-WIMP possibilities. The motivation behind non-WIMP phenomenology comes from the failure of standard WIMP in providing a suitable DM candidate. As described earlier, direct search experiments give the exclusion limit on the upper bound of DM-SM interaction. With the recent data and future predictions from such experiments, the allowed upper bound is gradually touching the neutrino floor, thus eliminating the prospect of standard WIMP models to account for the observed DM relic density. The search prospect of these models also become very limited because of the restricted DM-SM coupling, which is the only means of interaction for DM with SM. As a solution to this crisis, here we propose several non-WIMP scenarios

1. **Multi-component DM :** In multi-component DM models, the interplay of different DM components can successfully evade the difficulties faced by single component DM scenarios. Compared to standard WIMPs, DM-DM conversion is possible in MCDM, which can change DM number density without contributing to the direct detection cross-section. In search of such minimalistic setups accounting for a few GeV up to TeV mass range, we show that one needs to consider a two-component setup with an extended dark sector (contains more than one dark particles) in at least one of the components. Such enlarged dark sector often gives rise to rich collider phenomenology [51, 63, 151, 205]. We discuss two works on DM components of different spin and isospin in this context, with the indication of possibilities of collider probes.
2. **Coannihilating DM :** As discussed previously, the standard formalism of WIMP freeze-out is significantly modified if one considers almost degenerate dark sector particles, inducing coannihilation. Relic density achieved by coannihilation relaxes the parameter space in such scenarios, because coannihilation changes DM number density, but keeps the DM-SM interaction unchanged. Also, the compressed dark sector can be an interesting handle in the collider probe [73, 111]. On a slight change in parameters in such models, DM number changing channels can substantially vary and this, in turn, affects the collider phenomenology as well. We address such scenario in two works and show how dark leptons can lift the restrictions on scalar singlet WIMP and give rise to interesting DM production channels in the collider.
3. **Non-thermal DM :** We focus on the freeze-in of DM, where, as already explained above, the DM production is entirely different than the freeze-out of WIMP. But in minimalistic models, one can go from the non-thermal to the thermal regime with a proper variation of DM-SM as well as DM-DM interaction [172]. We propose a scenario where DM relic density can be achieved by freeze-in and freeze-out mechanisms, albeit in separate areas of parameter region. We also give a quantitative estimate of the limiting case where the transition between the thermal and non-thermal regime takes place. As a natural consequence of the model arrangement, one or more long-lived particles are produced, which have a delayed decay in the colliders, unlike any SM particle. These exotic phenomena have good potential to become the smoking gun for BSM physics at LHC soon enough [49].

We organize the thesis in three parts, each discussing one of the above non-WIMP scenarios. The chapters under each part discuss the separate papers. In the concluding part, we present the summary and highlights of the thesis. We highlight the significance and the impact of the included works in the context of frontier research in DM studies and particle physics phenomenology at large.

Part I

**Spinor/vector aided scalar multipartite
scenario**



In this part, we take up the possibility of multi-component DM scenario as a viable possibility where the current experimental bounds eliminate the parameter space of the conventional WIMP paradigm. Instead of a single DM candidate, as is the case for the conventional WIMPs, here we consider two WIMP DM candidates with possibilities of interaction between them. The interplay between the two DM candidates here is the key point of the two studies presented here. This is because this interaction is entirely independent of the coupling with the visible sector, and thus, through this coupling, relic density can be successfully explained despite keeping direct detection bounds below the observed limit. In the two chapters included here, we discuss minimal setups which feature the co-existence of two species of different spins and isospins as the two DM candidates. Detection prospect for such setups is also discussed, along with the limitations of the current search strategies and the possibilities of improvement in the near future with better sensitivity.

One of the DM components which is common in both the studies is the inert scalar doublet. The theory of this model is already discussed in Chapter 1. DM phenomenology of IDM has been elaborately studied in the literature [18, 33, 47, 76–78, 86, 88, 102, 104, 152, 188–190, 192–194, 204]. Apart from the dark matter phenomenology, collider searches from IDM are also performed extensively [25, 27, 33, 119, 146, 153, 157, 175, 197, 205]. IDM can provide a low mass dark matter candidate with mass smaller than 130 GeV and a high mass regime above 500 GeV [18, 204]. In the region in between, the inert dark matter is found to be underabundant due to the large annihilation into gauge bosons. However, this can be resolved if one assumes that dark matter is multicomponent in nature. Different multicomponent dark matter models including both thermal and non-thermal dark matter candidates have been explored in the literature [17, 28, 31, 56, 59, 61, 64, 65, 68, 71, 72, 93, 99, 126, 160, 177, 199, 201, 202, 222]. In multicomponent dark matter models one of the candidates shares a fraction of total DM relic abundance and the other candidate provides the required amount of relic density in order to satisfy the total DM relic density observed by Planck [16].



Interplay of scalar and fermionic components in a multicomponent dark matter scenario

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2.1 Motivation

In this chapter, we discuss a new scenario with one scalar and one fermionic DM particle coexisting to fulfil the relic density conditions, at the same time evading the direct detection possibilities so far. The key to our new proposal is the realization of the fact that the annihilation cross section has to be enhanced beyond what is provided by the Higgs portal interactions (which are constrained by the direct detection experiments), at the same time with a handle on the cross section provided by a tuneable coupling, unlike the case of annihilations enabled by the gauge couplings. This, in the proposed scenario, is achieved through the presence of a newly introduced fermionic field, which, along with the standard leptons share a Yukawa coupling with the dark-scalar doublet field. We shall show, that for sufficiently large ranges of the relevant Yukawa coupling and the mass of the new fermion, which are the two new parameters here, it is possible to have the desired relic density for the dark matter candidate. These new fermions should carry the same $\mathbb{Z}_2$ charge as that of the scalar DM, and therefore are required to be heavier than the DM itself. Thus, DM candidates with larger than about 100 GeV mass, require large values of the Yukawa couplings to compensate for the scaling down of the cross section with the correspondingly larger $\mathbb{Z}_2$ -odd fermions, which mediate the annihilations. Thus, it is still desirable to have non-singlet scalar fields to provide DM candidates. In addition to supporting to obtain the required relic density, the presence of such fermionic partners (considered here as electrically charged) provide a new handle to explore this scenario in collider experiments. However, the gauge mediated annihilations overkill the DM, as in the cases available in the literature. The presence of a fermionic partner and the additional channels of annihilation only worsen the situation. The introduction of another DM candidate not only salvages the situation but in a beautifully interlinked coexistence, provide sufficient relic density in a very large range of masses of both the DM candidates. We demonstrate this in the proposed model through the addition of a gauge singlet fermion field, stable under a different $\mathbb{Z}'_2$ symmetry. The fermionic DM interacts with the

SM particles through a singlet scalar portal, enabled by mixing of this neutral scalar with the SM Higgs field, leading to a natural way to explain direct detection limits through the smallness of the mixing without requiring fine-tuning of the parameters of the Lagrangian. The same portal coupling allows interaction with the scalar dark matter candidate field. Such portal interaction between the two dark matter fields enables the conversion of one type of dark matter to the other. In kinematically allowed phase space regions having a mass hierarchy with the fermionic dark matter heavier than the scalar one, a favourable condition arises with the help of such conversion processes. Specifically, this allows annihilation of the otherwise over-abundant fermionic DM to inject the otherwise under-abundant scalar DM to provide together with the desired relic density. This minimal two-component scalar-fermion dark matter scenario provides interesting additional features as a bonus, some of which are detailed in the rest of this chapter.

2.2 Model

The framework of the proposed model has the same gauge group as that of the SM. The particle content of the SM along with the Higgs doublet (Φ_1) is extended with the addition of a scalar doublet Φ_2 having hypercharge +1, two vector-like fermion singlets, χ and ψ with hypercharges -2 and 0 , respectively, and a singlet scalar field ϕ , with zero hypercharge. The new doublet field, Φ_2 and charged singlet fermion χ are considered odd under a discrete $\mathbb{Z}_2$ symmetry, while all other fields are considered even under this transformation. Similarly, the neutral singlet fermion ψ is taken to be odd under another $\mathbb{Z}'_2$ symmetry, while all other fields are considered even under this.

	Φ_1	Φ_2	χ_L, χ_R	ψ_L, ψ_R	ϕ
Y	+1	+1	-2	0	0
$\mathbb{Z}_2$	+	-	-	+	+
$\mathbb{Z}'_2$	+	+	+	-	+

Table 2.1: Relevant fields and their charges. Electromagnetic charges are given by $Q = t^3 + Y/2$.

With the above $\mathbb{Z}_2 \times \mathbb{Z}'_2$ discrete symmetry and the SM gauge symmetry, the new physics interaction part of the Lagrangian is given by

$$\begin{aligned} \mathcal{L} \supset & (D_\mu \Phi_2)^\dagger (D^\mu \Phi_2) + \bar{\chi} i \gamma^\mu D'_\mu \chi + \bar{\psi} i \gamma^\mu \partial_\mu \psi - M_\chi \bar{\chi} \chi - M_\psi \bar{\psi} \psi \\ & - (y_1 \bar{L} \Phi_2 \chi_R + h.c.) - y_2 \bar{\chi} \chi \phi - y_3 \bar{\psi} \psi \phi - V \end{aligned} \quad (2.1)$$

with the covariant derivatives $D^\mu = \partial^\mu + i g \tau \cdot W^\mu + i g' (Y/2) B^\mu$ and $D'^\mu = \partial^\mu + i g' (Y/2) B^\mu$, where g and g' are the corresponding gauge couplings and Y is the hypercharge. L denotes the SM lepton doublet field. The scalar potential is given by

$$\begin{aligned} V = & \mu_1^2 \Phi_1^\dagger \Phi_1 + \mu_2^2 \Phi_2^\dagger \Phi_2 + \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) \\ & + \lambda_4 |\Phi_1^\dagger \Phi_2|^2 + \frac{1}{2} [\lambda_5 (\Phi_1^\dagger \Phi_2)^2 + h.c.] + \mu_3^2 \phi^\dagger \phi + \lambda_6 (\phi^\dagger \phi)^2 \\ & + \frac{1}{2} [\mu_4 \phi (\Phi_1^\dagger \Phi_1) + \mu_5 \phi (\Phi_2^\dagger \Phi_2) + h.c.] + \frac{1}{2} [\mu_6 \phi^3 + \mu_7^3 \phi + \mu_8 (\phi^\dagger \phi) \phi + h.c.] \\ & + \lambda_7 (\phi^\dagger \phi) (\Phi_1^\dagger \Phi_1) + \lambda_8 (\phi^\dagger \phi) (\Phi_2^\dagger \Phi_2). \end{aligned} \quad (2.2)$$

With the standard Higgs field developing a vacuum expectation value (vev), $v = 246$ GeV, leading to the electroweak symmetry breaking (EWSB), the scalar fields may be expressed in the unitary gauge as

$$\Phi_1 = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} H^+ \\ \frac{H_0 + i A_0}{\sqrt{2}} \end{pmatrix}, \quad \phi = \frac{1}{\sqrt{2}} (h_s + i A_s). \quad (2.3)$$

To keep the $\mathbb{Z}_2$ symmetry intact, we disallow Φ_2 from developing a vev by setting $\mu_2^2 \geq 0$. Similarly, it is arranged so that ϕ does not generate a vev. The physical spectrum now has two charged scalars, $H^\pm$, one neutral scalar H^0 and a neutral pseudoscalar A^0 coming from Φ_2 , with the lightest of H_0 and A_0 becoming a dark matter candidate. We confine to the case of $M_{H_0} < M_{A_0}$. On the other hand, h and h_s mix to generate the two physical scalar bosons; the observed 125 GeV Higgs boson and another scalar boson denoted here by H and H_S , respectively. This mixing is parametrised with an angle α as

$$\begin{pmatrix} H \\ H_S \end{pmatrix} = \begin{pmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} h \\ h_s \end{pmatrix}. \quad (2.4)$$

The physical scalar masses are related to the quartic coupling λ_1 and the vev through the relation

$$M_{H_S}^2 - M_{A_S}^2 + M_H^2 = 2\lambda_1 v^2. \quad (2.5)$$

Considering λ_1 to be positive for the stability of the potential leads to the tree-level mass relation

$$M_{H_S}^2 + M_H^2 > M_{A_S}^2. \quad (2.6)$$

Mixing with the SM Higgs field allows the scalar component of the singlet field H_S to decay to the SM particles, thus allowing its mass to be practically unrestricted. On the other hand, the pseudoscalar component, A_s directly couples only to the new fermions and the inert Higgs field, and thus to allow tree-level decay its mass is required to be larger than twice the mass of the lightest dark matter candidate. The mass spectrum of the inert doublet field is not affected by other interactions, with the masses related to the parameters of the potential as in the pure IDM case given by [188]

$$\begin{aligned} M_{H^\pm}^2 &= \mu_2^2 + \frac{\lambda_3}{2} v^2 \\ M_{H_0}^2 &= M_{H^\pm}^2 + \frac{1}{2} (\lambda_4 + \lambda_5) v^2 \\ M_{A_0}^2 &= M_{H^\pm}^2 + \frac{1}{2} (\lambda_4 - \lambda_5) v^2. \end{aligned} \quad (2.7)$$

Besides, as explained in the introduction, the new scenario necessitates two additional charged leptons, and a neutral fermion ψ in the physical spectrum. The mass hierarchy of $M_{\chi^\pm} > M_{H_0}$ is maintained to allow the decay of $\mathbb{Z}_2$ odd fermion, $\chi^\pm$, whereas ψ , the fermionic component of the dark matter is made stable with the $\mathbb{Z}_2'$ symmetry. Apart from the mass relations in Eq. (2.7) above, the condition

$$\mu_7^3 = -\frac{\mu_4}{2} v^2 \quad (2.8)$$

is set to remove the linear term after the EWSB. This along with setting μ_3^2 , μ_6 , μ_8 , λ_6 and λ_7 to be positive definite makes sure that ϕ does not develop a non-zero vev. In our analysis, we have traded the parameter μ_3^2 for the physical mass, M_{A_S} , which are related through

$$\mu_3^2 = M_{A_S}^2 - \frac{\lambda_7 v^2}{2}. \quad (2.9)$$

In our choice of parameters, we have made sure that $\mu_3^2 \geq 0$, as required. We list the vertices and the corresponding Feynman rules relevant to the new degrees of freedom in Fig. 2.1. Further, we define

$$\lambda_L = \frac{1}{2} (\lambda_3 + \lambda_4 + \lambda_5) \quad (2.10)$$

as the combination of the couplings that are relevant to Higgs portal interaction of the scalar dark matter candidate involved in the direct detection experiments as well as the annihilation processes. Coming to the experimental constraints, LEP limits of

$$(M_{H_0} + M_{A_0}, 2M_{H^\pm}) > M_Z; \quad \text{and} \quad M_{H_0} + M_{H^\pm} > M_W \quad (2.11)$$

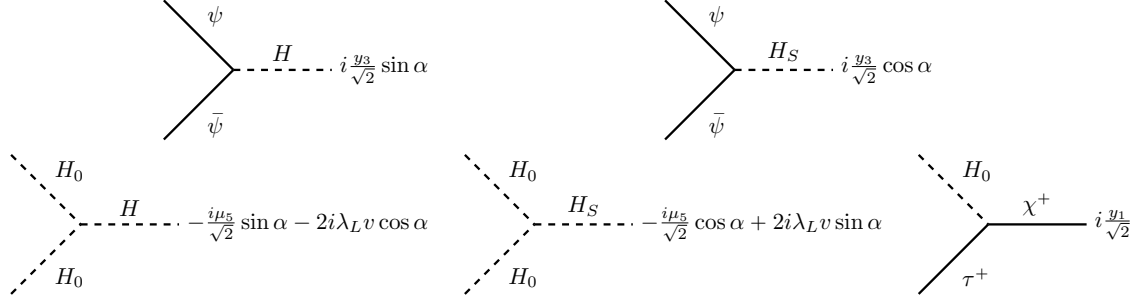


Figure 2.1: Feynman rules relevant to the annihilation of the DM candidates.

are obtained from non-observation of the decay of Z and W to the inert Higgs bosons. The precision electroweak measurements are sensitive to the mass splitting between the charged Higgs boson and the neutral ones, with the IDM contribution to the T -parameter given by

$$T_{IDM} = \frac{1.08}{v^2} (M_{H^\pm} - M_{H_0})(M_{H^\pm} - M_{A_0}). \quad (2.12)$$

The current experimental bound on the value of $T = 0.08 \pm 0.12$ [225] can be accommodated with at least one of the light neutral Higgs bosons having mass close to that of the charged Higgs boson.

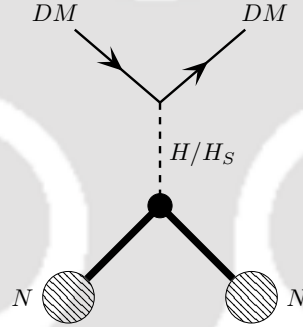


Figure 2.2: Schematic diagram showing the scattering of DM with nucleus considered in the direct detection experiments.

The direct detection depends on the elastic scattering of the dark matter candidate with the neutron and proton in the nucleus, which is mediated by the scalar bosons in our case, as shown in Fig. 2.2. This does not, therefore, distinguish whether the dark matter is a fermion or a scalar particle. While the scalar dark matter candidate has a direct coupling with the SM doublet field, the fermionic dark matter interacts with the visible sector only by the mixing between the new scalar field introduced and the standard Higgs boson. The scattering cross section of the scalar DM is dictated by the coupling λ_L in the IDM sector, and the newly introduced trilinear coupling μ_5 between ϕ and the inert doublet field Φ_2 . The process is mediated by H and H_S with the former case coupling to H_0 with appropriate combination of $\lambda_L v \cos \alpha$ and $\mu_5 \sin \alpha$, and the latter case with combination of $\lambda_L v \sin \alpha$ and $\mu_5 \cos \alpha$, as could be read from Fig. 2.1. The relevant spin-independent cross sections are given by

$$\begin{aligned} \sigma_{H_0 N \rightarrow H_0 N} &= \frac{f_N^2 M_N^4}{16\pi(M_{H_0} + M_N)^2} \left(\frac{\lambda_H}{M_H^2} \cos \alpha - \frac{\lambda_{H_S}}{M_{H_S}^2} \sin \alpha \right)^2, \\ \sigma_{\psi N \rightarrow \psi N} &= \frac{f_N^2 M_N^4}{16\pi(M_\psi + M_N)^2} \left(\frac{g_H}{M_H^2} \cos \alpha - \frac{g_{H_S}}{M_{H_S}^2} \sin \alpha \right)^2, \end{aligned} \quad (2.13)$$

where

$$\begin{aligned}
 \lambda_H &= -\frac{1}{\sqrt{2}} \left(\mu_5 \sin \alpha - 2\sqrt{2} \lambda_L v \cos \alpha \right), \\
 \lambda_{H_S} &= -\frac{1}{\sqrt{2}} \left(\mu_5 \cos \alpha + 2\sqrt{2} \lambda_L v \sin \alpha \right), \\
 g_H &= \frac{y_3}{\sqrt{2}} \sin \alpha, \\
 g_{H_S} &= \frac{y_3}{\sqrt{2}} \cos \alpha,
 \end{aligned} \tag{2.14}$$

and, f_N is the nuclear form factor. Notice that in the absence of mixing ($\sin \alpha = 0$), $\sigma_{H_0 N \rightarrow H_0 N}$ is reduced to the usual case in the pure IDM [47]. Further, $\sigma_{\psi N \rightarrow \psi N}$ is proportional to $(\sin \alpha \cos \alpha)^2$, reminding us that ψ interacts with the visible sector owing entirely to the mixing of ϕ with the doublet Higgs field.

The scattering cross-section of the DM components H_0 and ψ off the nucleon are sensitive to the couplings of the mediators with the DM, as well as the quarks in the nucleon. Also, the cross-section depends on the masses of DM and the mediators. For H_0 and ψ , the mediator couplings with DM are given in Eq. (2.14), whereas the mediator couplings with SM quarks are included in the cross-section expressions (Eq. (2.13)) in the form of the nuclear form factors and the mixing angles. Ultimately, apart from the DM mass, mediator mass (M_{H_S}) and the mixing angle, the DM-nucleon cross-sections are sensitive to μ_5 and λ_L for H_0 , as well as y_3 for ψ . The interplay of these parameters are shown in Fig. 2.3.

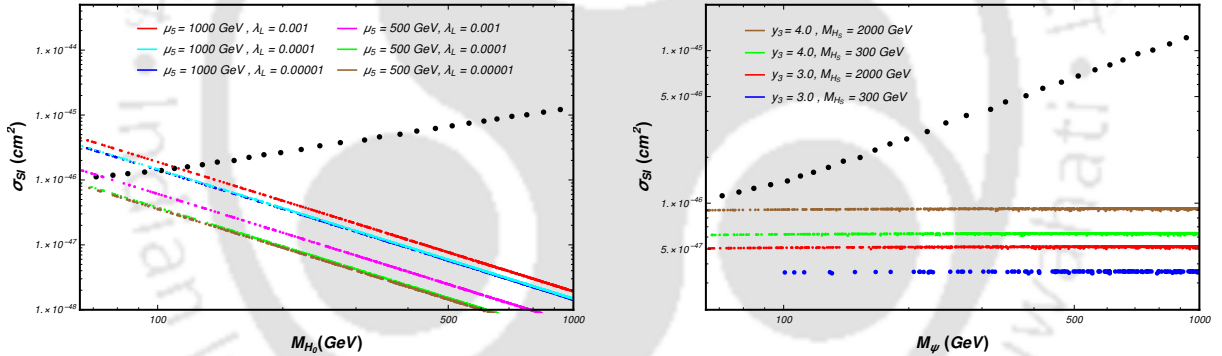


Figure 2.3: σ_{SI} vs M_{DM} for different couplings and mediator mass relevant to direction detection processes. The black dotted lines represent the current XENON IT bound [29].

To find the viable parameter values, we compare the direct detection cross section obtained using MICROMEGAS [50] with the XENON1T [29] bounds. In Fig. 2.3 (left) we plot the cross section for H_0 -nucleon scattering against M_{H_0} for two different values of $\mu_5 = 500$ GeV and 1000 GeV, with three different values of $\lambda_L = 10^{-3}, 10^{-4}, 10^{-5}$ in each case. As seen, the sensitivity of λ_L is insignificant for sizable μ_5 . While we have presented our results in Fig. 2.3 for slightly larger values of μ_5 , a similar pattern is seen in the case of smaller values ($\mu_5 \sim 100$ GeV) as well. Dependence on M_{H_S} is not presented, however, we have checked that it is not very significant. Notice that $\mu_5 \leq 500$ GeV is compatible with $M_{H_0} \geq 65$ GeV, and $\mu_5 \leq 1000$ GeV is compatible with $M_{H_0} \geq 100$ GeV. While considering the limits, we have not included the contribution of the other dark matter candidate, ψ , which means the scaling factor in $(\Omega_i/\Omega_{\text{tot}})\sigma_i$ [161] for the i^{th} DM component in multicomponent scenario is taken as 1. Moving on to the case of ψ (again, in the absence of H_0), cross section of ψ -nucleon scattering is plotted against M_ψ for different values of the $\psi\psi\phi$ coupling y_3 in Fig. 2.3 (right). Here again the mediators are H and H_S . However, the couplings are rather straight forward, unlike the previous case of H_0 -nucleon scattering, with the cross section proportional to $(y_3 \sin 2\alpha)^2$ irrespective the case with H or H_S as the propagator, as clear from Eq. (2.13) and discussion there. We have considered a fixed value of $\alpha = 0.0045$ and varied y_3 to see the effect of the coupling. Apart from the

coupling, the process depends strongly on M_{H_S} , the mediator mass. As can be seen from Fig. 2.3 (right), $y_3 = 4$ is allowed for the above value of mixing, and a rather heavy H_S with $M_{H_S} = 2$ TeV. Lighter H_S leads to more relaxed limit on the coupling, contrary to the naive expectation. This could be attributed to the destructive interference between the contributions from H_S mediation and H mediation. With these observations, we proceed to see the effect on the dark matter relic density.

In order to understand the compatibility of the model with the observations of relic density given by $\Omega h^2 = 0.1198 \pm 0.0012$ [16], where h denotes the Hubble parameter normalised to $100 \text{ km s}^{-1} \text{ Mpc}^{-1}$, we perform a relic density computation using the MICROMEAS package, scanning over the theoretically available parameter space regions. We consider two distinct scenarios with (i) $M_{H_0} \leq M_H/2$ and (ii) $M_{H_0} > M_H/2$. In the former case, the invisible decay of the Higgs boson to DM pairs will put restrictions on the couplings. Here, λ_L is the relevant coupling for $H \rightarrow H_0 H_0$ process, although, through the H - H_S mixing, this channel is also influenced by the trilinear coupling, μ_5 and the mixing angle, α . Considering the fermionic component of the DM, notice that the only interaction of ψ to other particles is facilitated by the singlet scalar field ϕ . Through mixing with the SM Higgs field, this leads to $\psi\psi H$ coupling of $y_3 \sin \alpha$. The present LHC bound on $H \rightarrow \text{invisible}$ decay width is restricted to about 20% [21], leading to a constraint on $y_3 \leq 0.02$ for the maximum allowed $\alpha \sim 0.33$. In our analysis, in the region $M_{H_0} \leq M_H/2$ and $M_\psi \leq M_H/2$, we have discussed two scenarios. One with $\lambda_L = \mu_5 = y_3 = 0$ as the most conservative approach, so that the invisible decay of H to the dark matter particles is disallowed, and the second case in which we relax this with non-zero values of these couplings, which are compatible with the present experimental limit on the invisible decay of Higgs boson. However, in the first case setting y_3 to be zero whenever $M_\psi \leq M_H/2$ results in an overabundance of ψ and consequently ruling out these mass ranges for ψ . Therefore, in those discussions we shall consider $M_\psi > M_H/2$. Finally, the Yukawa coupling y_1 between the SM leptons and dark fermion χ with the inert scalar field, when allowed for the first two generations may induce larger than desired $(g - 2)$ for electron and muon. Besides, the presence of non-diagonal couplings can induce undesired lepton-flavour-violating processes like $\mu \rightarrow e\gamma$. To avoid such effects, we consider the couplings to be diagonal and χ that couples to the first two generations to be very heavy, along with suitably chosen small values of the corresponding couplings. Therefore, in our further discussion we consider only the third generation coupling to be present, leading to $\chi\tau H_0$ interaction. As is clear from the discussion below, this coupling and the corresponding interaction makes significant impact only in the light H_0 case, i.e., $45 \text{ GeV} \leq M_{H_0} \leq 80 \text{ GeV}$. Before embarking on our numerical analysis, we shall look at the details of the above scenarios.

2.2.1 Scenario 1 : $M_{H_0} \leq M_H/2$

Keeping all other scalar masses larger than its mass, H_0 annihilates into the SM leptons and light quarks, mediated through H and H_S in the s -channel, and through the newly introduced heavy fermion, χ in the t -channel. Notice that, in the absence of χ the annihilation is enabled through the s -channel process, $H_0 H_0 \rightarrow (H, H_S) \rightarrow f\bar{f}$, where f is the SM fermion. The cross section of this process is proportional to

$$\left(\frac{\lambda_H}{M_H^2} \cos \alpha - \frac{\lambda_{H_S}}{M_{H_S}^2} \sin \alpha \right)^2,$$

the same coupling factor appearing in the direct detection cross section in Eq. (2.13), and, therefore, stringently constrained. In that case, the cross section is not sufficient to bring down the relic density to the observed value. Thus, the presence of an additional scalar singlet does not help improve the situation. The presence of χ , however, changes the scenario by adding to the cross section a t -channel process $H_0 H_0 \rightarrow \tau\tau$, mediated by χ . The corresponding Feynman diagram is given in Fig. 2.4a), and the cross section is given in Eq. (2.A.1).

This additional cross section can be tuned with the help of the unrestricted $H_0\chi\tau$ Yukawa coupling y_1 to get the required relic density. On the other hand, for large values of y_1 the annihilation of H_0 can make it underabundant. With $M_\psi > M_{H_0}$, the contribution to relic density from the H_0 , denoted

by $\Omega_1 h^2$, is independent of M_ψ itself, except for a small dependence on the conversion process $\psi\psi \rightarrow H_0 H_0$. This additional creation of H_0 is controlled by the trilinear and quartic couplings μ_5 and λ_L , the mass of H_S , and the Yukawa coupling, y_3 , as is clear from the Feynman diagram in Fig. 2.4b). The cross section for this process is given in Eq. (2.A.2). Setting $\lambda_L = \mu_5 = 0$ to avoid invisible decay of H , as stated above, makes this channel irrelevant. In addition to $\Omega_1 h^2$, the total relic density ($\Omega_{\text{tot}} h^2$) has the fermionic component, $\Omega_2 h^2$, so that $\Omega_{\text{tot}} h^2 = \Omega_1 h^2 + \Omega_2 h^2$. Note that $\Omega_2 h^2$ is controlled by the $\psi\psi$ annihilation into the SM states. Before the gauge boson annihilation channels open up for ψ at $M_\psi \sim 80$ GeV, the fermionic component $\Omega_2 h^2$ is larger than the allowed relic density, unless the other annihilation channel $\psi\psi \rightarrow ff$, where f denotes the SM fermions, is sizable. This latter process, mediated by the singlet component of H and H_S depends on the combination of $y_3 \sin 2\alpha$ through the $(\psi\psi H, ffH)$ pair of interactions (see Fig. 2.1) in the H mediated case, and through $(\psi\psi H_S, ffH_S)$ pair in the H_S mediated case. The Feynman diagram for this process is given in Fig. 2.4c), and cross section in Eq. (2.A.5). On the other hand, for $M_\psi \geq 80$ GeV, the $\psi\psi \rightarrow VV$, where $V = W, Z$, the channel allows a considerable reduction in $\Omega_2 h^2$, opening large parameter space region compatible with the current measurements. The cross section for these processes is given in Eqs. (2.A.3) and (2.A.4). While these processes are also suppressed by the Higgs mixing, these are much more significant compared to the $f\bar{f}$ annihilation channel with

$$\frac{\sigma_{\psi\bar{\psi} \rightarrow VV}}{\sigma_{\psi\bar{\psi} \rightarrow f\bar{f}}} = \frac{1}{x N_C} \frac{M_V^4}{4M_\psi^2 m_f^2} \frac{\beta_V}{\beta_f} \left(3 - \frac{4M_\psi^2}{M_V^2} + \frac{16M_\psi^4}{M_V^4} \right), \quad (2.15)$$

where $x = 8$ for $V = W$ and $x = 16$ for $V = Z$, and the color factor, $N_C = 1$ for leptons and $N_C = 3$ for quarks. As these gauge annihilation channels are s -channel processes mediated by H and H_S , the allowed parameter region is expected to be around the resonant condition $M_{H_S} \sim 2M_\psi$. Away from this region, for $M_\psi \geq M_{H_S}$, the possibility of ψ annihilation is controlled by the $\psi\psi \rightarrow H_S H_S$ process. While the H/H_S mediated s -channel depends on the trilinear couplings μ_6 and μ_8 , the more important t -channel (see Fig. 2.7b) for the Feynman diagram) depends on the Yukawa coupling y_3 . The cross section for this process is given in Eq. (2.A.6). There is also a less relevant $\psi\psi \rightarrow HH$ possibility, which, however, is proportional to $(y_3 \sin \alpha)^2$.

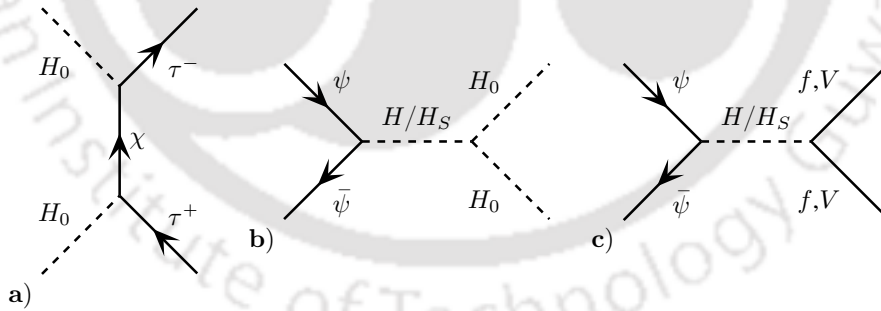


Figure 2.4: Feynman Diagram showing (a) the t -channel annihilation channel of H_0 , (b) the conversion of ψ to the scalar DM candidate and (c) the annihilation of ψ into SM final states, where $V = W, Z$

Interplay of the DM components: Note that, the individual components should be in an underabundant state so that the total relic density is within the desired bound. A situation that warrant particular attention is the case of $M_H/2 < M_\psi < 80$ GeV. Here, the annihilation of ψ is decided mainly by two processes leading to ff and $H_0 H_0$ final states. Also, when $M_\chi < M_\psi$ it is possible to have a pair of ψ annihilating into a pair of χ . Feynman diagram for this process is given in Fig. 2.7a), and the corresponding cross section is given in Eq. (2.A.8). Since the conversion channel depends on the couplings which are constrained by the direct detection experiments and the invisible Higgs decay bounds, the presence of χ is necessary to have the combined relic density in the allowed limit. For non-zero values

of λ_L , μ_5 and y_3 , the only substantial change is in $\Omega_2 h^2$, which could now be reduced to within the observed bound even for $M_\psi < M_V$ due to the additional H and H_S mediated channels.

2.2.2 Scenario 2: $M_{H_0} > M_H/2$

In this region λ_L and μ_5 are relatively unconstrained, opening up possibilities beyond what is discussed in *Scenario 1* with $M_{H_0} \leq M_H/2$. This along with the fact that H_0 can be lighter than ψ makes the Boltzmann's equations for each of these species interdependent. Thus, $\Omega_1 h^2$ would now depend not only on the mass and couplings of ψ but also on its number density. The process $\psi\psi \rightarrow H_0 H_0$ is now relevant to both $\Omega_1 h^2$ and $\Omega_2 h^2$, dictated by μ_5 , λ_L , y_3 and the mixing angle α . It may be noted, that with the only inert scalar present, a substantial region of parameter space with $80 \lesssim M_{H_0} \lesssim 500$ GeV is underabundant due to the large annihilation into the gauge bosons. Now with the presence of ψ and the possibilities mentioned above, it opens up a large window of DM mass region accessible.

Indirect search limits: The parameters which are sensitive to couplings of both DM components to SM are suppressed to very low values ($\lambda_L \sim \mathcal{O}(10^{-4})$, $\sin \alpha \sim \mathcal{O}(10^{-4})$) in order to evade direct search constraints and highlight the effect of DM-DM conversion in the relic density. This makes the indirect search prospects insignificant. However, the scalar DM component, being inert doublet DM, has gauge interactions which is not suppressed. We discussed this in Section 3.5 and showed that $\langle \sigma v \rangle$ is below the bounds from Fermi-LAT and HESS for couplings ~ 2 orders of magnitude larger than what is discussed here.

2.3 Numerical Results

We now come to our numerical results in this section. The quartic couplings, λ_2 and λ_6 involve four-point self-interactions and therefore do not influence the relic density computations. Similarly, λ_7 , μ_4 , μ_6 and μ_8 would influence the annihilation through s -channel into singlet scalar or the SM Higgs boson mediated by these same scalar fields. λ_8 , which can involve in the $H_0 H_0$ annihilation to the singlet scalars through a four-point interaction on the other hand will not have much influence on the analysis. We have, therefore, fixed $\lambda_2 = \lambda_6 = \lambda_7 = \lambda_8 = 0.1$, and $\mu_4 = \mu_6 = \mu_8 = 100$ GeV in our study. This leaves the Yukawa couplings, y_1 , y_2 , y_3 , the quartic coupling combination λ_L , and the trilinear coupling between the scalar singlet and the inert doublet, μ_5 , apart from the relevant masses, which we consider as independent parameters in our numerical analysis.

2.3.1 Scenario 1: $M_{H_0} \leq M_H/2$

As discussed in Section 2.2.1, in Case 1 we set $\lambda_L = 0$ and $\mu_5 = 0$ and $M_\psi > \frac{M_H}{2}$ in this scenario and in Case 2 we consider non zero values of these parameters compatible with both invisible Higgs decay width and direct detection bounds. Considering the LEP constraint, we keep $M_{H_0} > 45$ GeV. Other parameters are considered as given in Table 2.2. This choice corresponds to the quartic couplings of the Lagrangian in the ranges, $\lambda_1 = (0.26, 0.56)$, $\lambda_4 = (0.262, 0.335)$, $\lambda_5 = (-0.377, -0.297)$, $\lambda_3 = (0.036, 0.042)$.

In the case of IDM Higgs bosons, LEP rules out the region where $M_{H_0} < 80$ GeV, $M_{A_0} < 100$ GeV and $M_{A_0} - M_{H_0} > 8$ GeV, since it would lead to visible di-lepton or di-jet signals [192]. At the same time, mass splitting below 8 GeV does not support relic bound [51]. We have checked that the situation does not change in the present model. Further, LEP-II constrains $M_{H^\pm} > 70$ GeV from non-observation of $e^+e^- \rightarrow H^+H^-$ production [203]. The Electroweak Precision Measurements (EWPM) require product of the mass splittings, $(M_{H^+} - M_{A_0})(M_{H^+} - M_{H_0})$ to be small [189]. These considerations have led to deciding M_{A_0} to be larger than 100 GeV for the range of M_{H_0} considered here, keeping M_{H^+} close to M_{A_0} . At the same time, $M_{A_0} > 180$ GeV in this set up would correspond to $|\lambda_{4,5}| > 1$. We, therefore, set a mass splitting of $M_{A_0} - M_{H_0} = 60$ GeV as our conservative choice. The choice

M_{H_0}	$45 \leq M_{H_0} \leq 63$			
M_{A_0}	$M_{H_0} + 60$			
$M_{H^\pm}$	$M_{H_0} + 65$			
M_{A_S}	$121 \leq M_{A_S} \leq 500$			
M_{H_S}	$M_{A_S}, M_{A_S} + 50$			
M_χ	(75, 130, 200)			
y_1, y_2, y_3	$0 \leq y_i \leq 3$			
μ_4, μ_6, μ_8	100			
$\lambda_2, \lambda_6, \lambda_7, \lambda_8$	0.1			
α	0.0045			

Case 1		Case 2
$65 \leq M_\psi \leq 1000$	M_ψ	$10 \leq M_\psi \leq 1000$
0	λ_L	0.0001
0	μ_5	100

Table 2.2: Parameters considered for the DM relic density study in scenario 1. All mass parameters are in GeV.

of $M_{A_s} > 2M_{H_0}$ is made to allow tree-level decay of A_s . As indicated by Eq. (2.5) we shall keep the mass splitting between that of H_S and A_s small enough to keep λ_1 small. At the same time, to keep λ_1 positive all through the parameter region, we make the conservative choice of $M_{H_S} > M_{A_s}$.

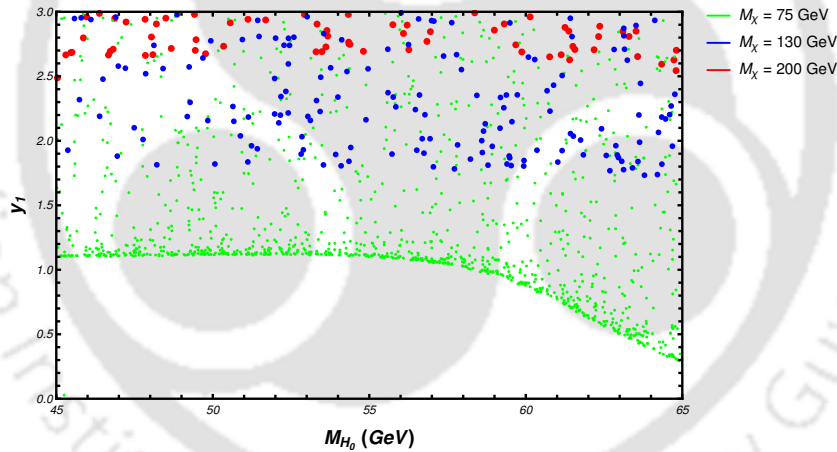


Figure 2.5: Region of $y_1 - M_{H_0}$ plane compatible with total relic density within the allowed range for selected values of M_χ . Other variables are varied as per Table 2.2 for Case 1.

With the above choice of parameters, micrOMEGAs is used to perform a random scan to compute the relic density to find compatible regions. With $\lambda_L = 0 = \mu_5$ disabling all the H and H_S mediated processes, the cross section is dominated by the χ mediated processes discussed in the previous section. The compatible regions in $y_1 - M_{H_0}$ plane for three different illustrative choices of M_χ is presented in Fig. 2.5. As expected, heavier mediator requires larger coupling for the same level of the cross section to satisfy the bounds. The other candidate for dark matter, ψ couples directly only to the singlet scalar field, ϕ , as elaborated in the previous section. Thus, its annihilation process is mediated by H_S , whose major component is ϕ , and the observed 125 GeV resonance of H having a small admixture of ϕ enabled by the nonzero mixing angle, α . The number density of otherwise overabundant ψ is reduced with the opening of the gauge annihilation channel with $M_\psi \geq M_W$ as clear from the top-right inset of Fig. 2.6 in the $\Omega_2 h^2$ vs M_ψ plot. For Case 2 in Table 2.2 of parameter choice where λ_L and μ_5 are non-zero, the only visible change in the results is shown in the top left-inset. Here $\Omega_2 h^2$ is reduced to the observed bound from overabundance even for $M_\psi < M_W$, thanks to the additional annihilation of ψ into SM

leptons and H_0 now made possible with the non-zero couplings. Since all these channels are H and H_S mediated, we also see the s-channel resonance effect at $M_\psi = M_H/2$ in this plot. H_S dependence in the

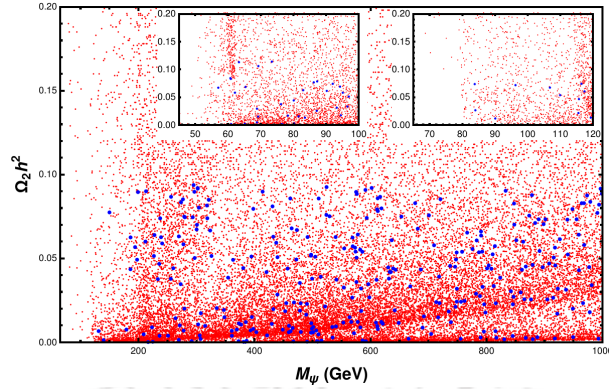


Figure 2.6: $\Omega_2 h^2$ vs M_ψ for the parameter region considered in Table 2.2. Blue points are a subset satisfying the total relic density $\Omega_{\text{tot}} h^2$ within the allowed range. Small mass region is enlarged for clarity : *Top-left inset*: for parameter region Case 2, *Top-right inset*: for parameter region Case 1 in Table 2.2

relic density calculations come through its mediation of ψ annihilation processes as well as annihilation of ψ pairs into H_S pairs. Thus the resonant enhancement of the cross section indicate that the compatible region has $M_{H_S} \sim 2M_\psi$, as clear in Fig. 2.7. When the full range of M_ψ is considered, the annihilation channels of $\psi\psi \rightarrow H_S H_S$, HH , $\chi^\pm \chi^\mp$ open up. The mass relation between M_ψ and M_{H_S} is no longer linear due to these new channels, leading to the scattered points in the high mass range starting from 125 GeV. Notice that $\psi\psi \rightarrow H_S H_S$ t -channel process and $\psi\psi \rightarrow \chi\chi$ process have negligible dependence on the mixing angle α , whereas for all other processes $\sin \alpha$ appears in combination with y_3 . Thus, in most situations, a change in α is compensated by a corresponding change in y_3 .

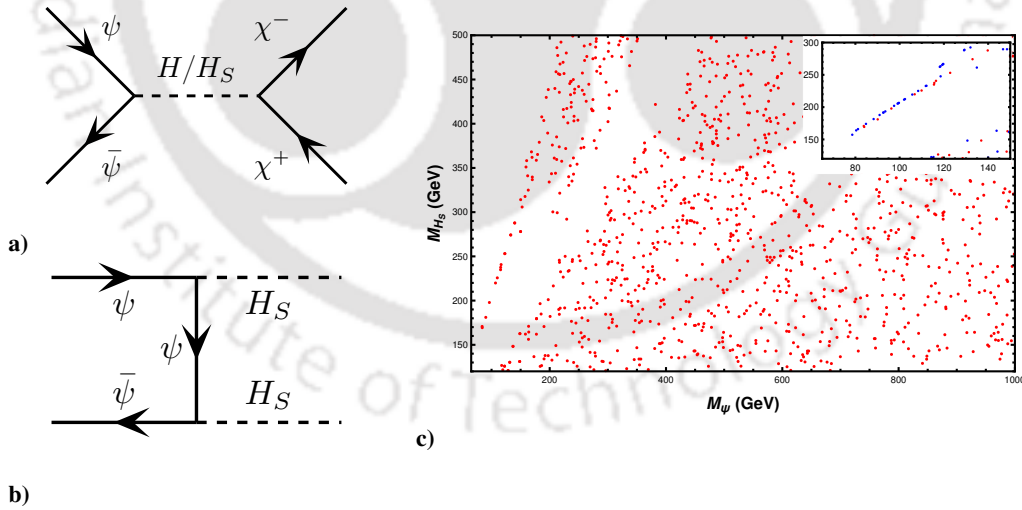


Figure 2.7: (Right): Regions of $M_{H_S} - M_\psi$ with $\Omega_{\text{tot}} h^2$ within the allowed relic density range. *Inset*: Low mass region is enlarged for clarity. The blue points correspond to $\Omega_2 h^2$ alone within the experimental bound. (Left): Feynman diagrams showing the relevant annihilation channels.

For $M_\psi \geq M_\chi$ the s -channel ψ annihilating into χ process is mediated by the singlet scalar and is proportional to the product of the Yukawa couplings, $y_2 \cdot y_3$. In Fig. 2.8 regions on $y_3 - M_\psi$ plane compatible with the total relic density bound is shown for $M_\chi = 100$ GeV, 150 GeV and 200 GeV, in each case for two different choices of y_2 values. $M_\psi = M_\chi$ threshold is seen in all the cases considered. Further, for small values of M_χ , the contribution from $\Omega_1 h^2$ to the relic density is negligible, as the χ mediated t -channel cross section being large washes it out, minimising the spread in the allowed

region. However, for $M_\chi = 200$ GeV and above, $\Omega_1 h^2$ has a non-negligible contribution, as clearly indicated in the larger area of allowed regions. The dependence on y_2 compared to y_3 is somewhat trivial as mentioned above. Points with different y_2 values shown unambiguously bring out the role of the $\psi\psi \rightarrow \chi\chi$ process. The split lines in the case of smaller y_2 values are due to the effects of other processes like ψ pair annihilating into the singlet scalars.

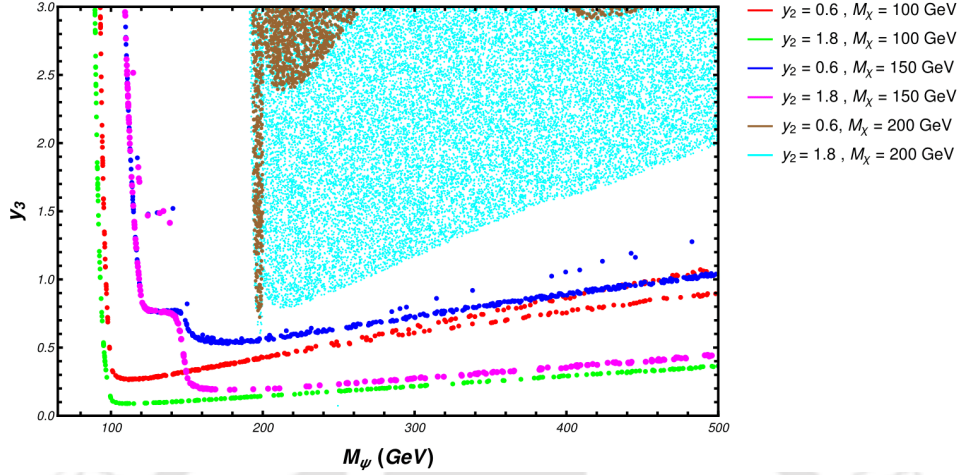


Figure 2.8: Regions of $y_3 - M_\psi$ plane compatible with $\Omega_{\text{tot}} h^2$ satisfying the experimental bounds, for specific values of y_2 and M_χ .

2.3.2 Scenario 2: $M_H/2 \leq M_{H_0}$

M_ψ	$65 \leq M_\psi \leq 1000$	M_χ	1200
M_{H_0}	$65 \leq M_{H_0} \leq 1000$	y_1, y_2, y_3	$0 \leq y_i \leq 3$
M_{A_0}	$M_{H_0} + (20, 60, 1)$	λ_L	0.0001
$M_{H^\pm}$	$M_{H_0} + (20, 65, 1)$	μ_5	(0, 100, 200, 500)
M_{A_S}	$131 \leq M_{A_S} \leq 1000$	μ_4, μ_6, μ_8	100
M_{H_S}	$M_{A_S} + 50$	$\lambda_2, \lambda_6, \lambda_7, \lambda_8$	0.1
		α	0.0045

Table 2.3: Parameters considered for the DM relic density study in scenario 2. All mass parameters are in GeV.

With non-zero values of λ_L and μ_5 , corresponding to the couplings of $\psi\psi\phi$ and $H_0 H_0 \phi$, respectively, the process $\psi\psi \rightarrow H_0 H_0$ mediated by ϕ (Fig. 2.4b)) makes the interaction between the two dark matter components more relevant. The set of coupled Boltzmann's equations make $\Omega_1 h^2$ dependent directly on $\Omega_2 h^2$ and vice versa. The value of $HH_0 H_0$ coupling is kept at $\lambda_L = 10^{-4}$, to respect the direct detection limit. Four different illustrative values including zero are considered for μ_5 .

The first row of Fig. 2.9 shows the variation of $\Omega_1 h^2$ vs M_{H_0} for specific choice of μ_5 . In the case of $\Omega_1 h^2$ with $\mu_5 = 0$ dark matter annihilating into SM particles alone is possible, like in the pure IDM case. Consequently, for $M_{H_0} \geq M_W$ the annihilation cross section controlled by gauge coupling leaves H_0 underabundant. However, with non-zero value of μ_5 in the region $M_\psi \geq M_{H_0}$, $\psi\psi \rightarrow H_0 H_0$ boosts up the relic density. This is illustrated in the top-middle plot of Fig. 2.9. For very large μ_5 , increase in annihilation cross section reduces the relic density below the observed limit, making the model less

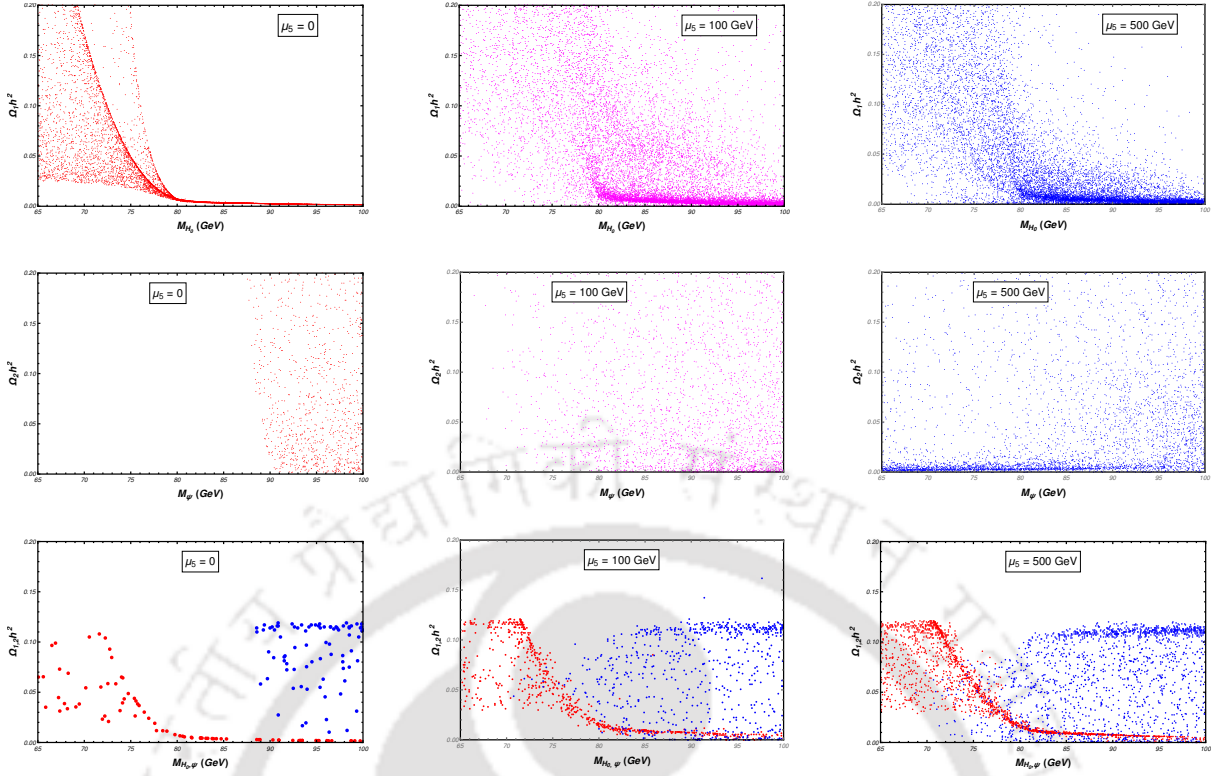


Figure 2.9: In *Top Row* $\Omega_1 h^2$ as a function of M_{H_0} (GeV) and in *Middle Row* $\Omega_2 h^2$ as a function of M_ψ (GeV) for illustrative values of μ_5 (GeV), indicating the importance of the conversion channel ($\psi\psi \rightarrow H_0 H_0$). The total number of points generated are the same in all the cases. *Third Row* shows selected points from the corresponding plots in the top and middle rows compatible with $\Omega_{\text{tot}} h^2$ within the observed bound. *Red* points correspond to $\Omega_1 h^2$ vs. M_{H_0} , while *blue* points correspond to $\Omega_2 h^2$ vs. M_ψ .

compatible as indicated in the top-right plot of Fig. 2.9. Second row of Fig. 2.9 shows the variation of $\Omega_2 h^2$ vs. M_ψ . In this case, ψ is overabundant until the gauge boson annihilation channels open up around $M_\psi \sim M_W$ when μ_5 is set to zero, as shown in the left plot. On the other hand, with $\mu_5 \neq 0$, the additional channels including the annihilation into $H_0 H_0$, $A_0 A_0$ and $H^+ H^-$ bring down the relic density to acceptable levels for $M_\psi \leq M_W$ as well.

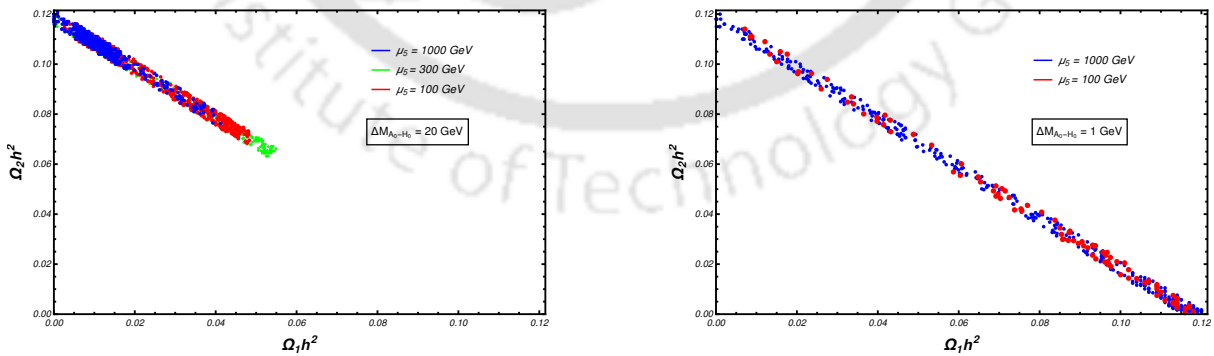


Figure 2.10: Correlation between $\Omega_1 h^2$ and $\Omega_2 h^2$ for points satisfying total relic density bound for illustrative values of μ_5 . Two different mass splitting of $\Delta M_{A_0-H_0} = 1$ GeV (*right*) and 20 GeV (*left*) of the IDM neutral Higgs bosons show the importance of co-annihilation channels.

Fig. 2.10 shows the correlation between $\Omega_1 h^2$ and $\Omega_2 h^2$ at a mass splitting $\Delta M_{H_0-A_0}=1$ GeV (*right*) and 20 GeV (*left*) for a mass range of $500 \leq M_{DM} \leq 1000$ GeV. When the mass splitting is small (of the order of 1 GeV), co-annihilation between the inert scalars counters the gauge suppression and increases the scalar relic density substantially. In this case, the effect of μ_5 is negligible as seen from

Fig. 2.10 (right). On the other hand, for a larger mass splitting, when the co-annihilation is suppressed, one may achieve significant contribution from $\Omega_1 h^2$ for a suitably chosen value of μ_5 . We find that the best-case scenario corresponds to a value of $\mu_5 \sim 300$ GeV. Hence in the multicomponent scenario, the contribution from $\Omega_1 h^2$ is boosted up compared to the single component case, thanks to the conversion from the fermionic component. We would like to reiterate the advantage of the multicomponent case considered here, which deviate from the purely IDM like scenario, where this mass range of DM is available only for a closely degenerate case of $M_{H_0} \sim M_{A_0}$.

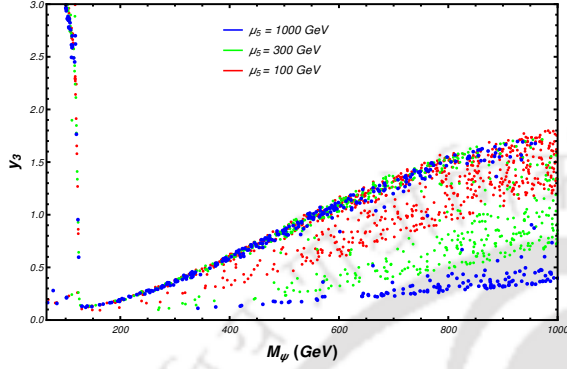


Figure 2.11: Regions of $y_3 - M_\psi$ plane compatible with the relic density bound for specific choices of μ_5 values, indicating the strong presence of the conversion channel $\psi\psi \rightarrow H_0 H_0$ in the resonant region with $M_{H_S} = 2M_\psi$.

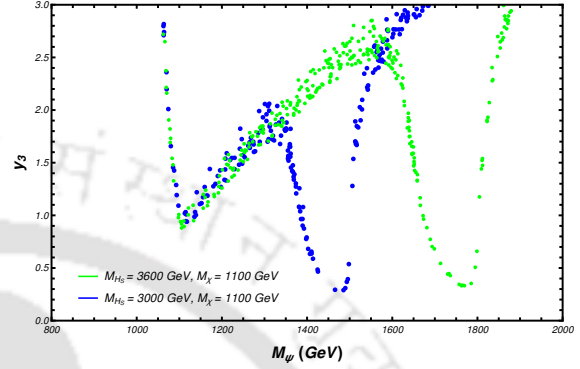


Figure 2.12: y_3 vs M_ψ for fixed M_{H_S} and M_χ values, showing the presence of the annihilation channel $\psi\psi \rightarrow \chi\chi$ and the resonant behaviour with $M_{H_S} = 2M_\psi$.

Presenting the correlation between M_ψ and the Yukawa coupling y_3 between $\psi\psi H_S$ Fig. 2.11 shows that higher value of μ_5 corresponds to lower y_3 for a fixed value of M_ψ , implying that $\psi\psi \rightarrow H_0 H_0$ is a dominant channel as per relic density consideration. Remember that μ_5 corresponds to the $H_S H_0 H_0$ interaction. To maintain the cross section at a certain value to follow the relic bound, higher μ_5 will correspond to lower y_3 and vice versa. All points here satisfy the total relic density bound. Notice that we have considered the resonance condition $M_{H_S} = 2M_\psi$, while M_ψ varies in the full range of Scenario 2, bringing in the relevant s -channel annihilation of ψ mediated by H_S .

In the large M_ψ region, the DM annihilation into χ pair opens up adding further possibilities. The $\psi\psi \rightarrow \chi\chi$ is an s -channel process mediated via H_S . Thus, the couplings y_2 and y_3 along with the masses of H_S , χ and ψ decide the cross section. Fig. 2.12 shows the allowed regions of $y_3 - M_\psi$ plane for specifically chosen values of the Yukawa coupling y_2 and the masses M_χ and M_{H_S} , which is set to $y_2 = 0.5$. The threshold is seen as $M_\psi \sim M_\chi$, which is set to 1100 GeV here. Further, the depression around the resonant condition of $M_{H_S} = 2M_\psi$ is seen, as expected in the s -channel process. Since in this region, the mass splitting between H_0 and A_0 is kept at 10 GeV, $\Omega_1 h^2$ is always small, hence the dominant contribution in $\Omega_{\text{tot}} h^2$ comes from $\Omega_2 h^2$.

2.4 Search prospects at the colliders

Since the additional DM component interacts with the SM only through the Higgs portal, there is no scope to distinguish this scenario from other model consisting of another DM particle of different spin interacting similarly (as discussed in the next chapter). However, some distinguishing features may be worth looking for. In the minimal scenario, the extra dark fermion must require a singlet scalar S as a portal to SM, as discussed in this chapter. But for another DM particle of different spin as discussed in the next chapter for vector case, this is not required. Therefore, for the extra fermion case, there is Higgs as well as Higgs-like scalar portal whereas for the other case, only the Higgs portal exists. At the LHC, the DM signal might be difficult to probe owing to the electroweak nature of the new physics sector, and hence the production of them through quark or gluon initiated processes being suppressed. However, at the ILC, one can probe via $e^+e^- \rightarrow ZH(S)$ channel as shown in the diagram below. Because of the

presence (absence) of $e^+e^- \rightarrow ZS$ channel for extra fermion (vector) dark matter in the dilepton+MET or dijet+MET signal processes, one can distinguish the two models from the peak and tail positions of suitable distributions. This can be relevant for dark matter of mass > 100 GeV, because, DD limits allow the Higgs portal coupling of order $\sim \mathcal{O}(10^{-2})$, which make the cross-sections significant in the resonance region where $m_S > 2m_{DM}$, m_S being the extra scalar mass.

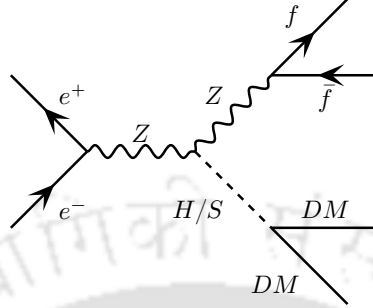


Figure 2.13: Possible production channel of DM at the ILC

2.5 Summary

The scalar-fermionic multipartite scenarios discussed in the literature have the DM candidates as gauge singlets and consequently highly constrained by the direct detection experiments and relic density limits. In these scenarios, the direct detection prefers small portal couplings, which consequently provide much smaller cross section than required to contain the overabundance of the dark matter relic density. On the other hand, the gauge non-singlet scalar dark matter models like the IDM overkills the dark matter leading to underabundance except when they are close to a TeV of mass. We have demonstrated that the presence of a fermionic partner (denoted as χ) to the inert scalar would alleviate these difficulties, opening up the low mass regions as well. Along this line, we present a novel scenario with the possibility of scalar and fermionic dark matter coexisting, compatible with all the experimental bounds including the relic density measurements, the direct detection limits and the collider constraints. We consider the presence of a gauge singlet fermion interacting with the visible sector through Higgs portal couplings, along with the gauge doublet scalar dark matter. A wide range of parameter space (10 GeV - 2 TeV) for dark matter mass is considered and the possible signatures are analyzed. We find that the entire mass range is compatible with the relic density and direct detection bounds. The dark matter particles interact among themselves opening possibilities of conversion from one type to the other leading to interesting phenomenology and compensate for the underabundance of the individual relic density of H_0 in the otherwise not compatible range of $80 \text{ GeV} \lesssim M_{H_0} \lesssim 500 \text{ GeV}$. For the entire mass range of the inert scalar H_0 , the fermionic dark matter candidate ψ contributes to the observed relic density starting from a few GeV to the TeV range. The lepton portal annihilation channels contribute to the relic density of H_0 , denoted by $\Omega_1 h^2$, without adding to the direct detection cross-section, being a t -channel processes mediated by the fermionic partner χ . Hence keeping λ_L fixed at an admissible low value at par with the direct detection limits, the lepton portal couplings and mass of χ can be adjusted to get the correct relic density for H_0 . Owing to the conversion of fermionic dark matter pair into a scalar dark matter pair substantial contribution of $\Omega_1 h^2$ in the total relic density is possible. The fermionic component can suffice for the deficit in the total relic density as well. On the other hand, for large rate of the annihilation of H_0 into SM through the Higgs portal channels, $\Omega_1 h^2$ could become very small. In such scenarios, the fermionic component dominates the scene with the dark matter scenario effectively becoming a single-component case. In the large mass region (500 GeV - 2 TeV) typical IDM contributes substantially to the relic density for very low mass splitting between the inert scalars, thanks to the now relevant co-annihilation channels. However, the effect of co-annihilation is negligible with larger mass splitting

making it non-compatible with the relic density measurements. In the model discussed here, fermionic to scalar dark matter conversion permits even a larger mass splitting to produce the correct relic density.

Finally, we expect that the model can bring in interesting collider phenomenology with the fermionic partner, χ of the inert scalar doublet changing the production and decay patterns of the IDM charged scalars in the mass range that could be probed at LHC. We defer a detailed collider study for future work.

Appendix

2.A Expressions for Cross sections

Cross section for $H_0 H_0 \rightarrow \tau^+ \tau^-$, neglecting the τ mass, is given by

$$\sigma_{H_0 H_0 \rightarrow \tau^+ \tau^-} = \frac{1}{16\pi s^2} \frac{1}{\beta_{H_0}^2} \int_{t_{min}}^{t_{max}} |\overline{\mathcal{M}_{H_0}^2}| dt \quad (2.A.1)$$

where $t_{min} = -(s/4)(1 + \beta_{H_0})^2$, $t_{max} = -(s/4)(1 - \beta_{H_0})^2$ with $\beta_{H_0} = \sqrt{1 - 4M_{H_0}^2/s}$, and the square of the invariant amplitude,

$$|\overline{\mathcal{M}_{H_0}^2}| = \frac{y_1^4}{4} [M_{H_0}^4 + t(s - M_{H_0}^2) + t^2] \left\{ \frac{2}{(u - M_\chi^2)(t - M_\chi^2)} - \frac{1}{(u - M_\chi^2)^2} - \frac{1}{(t - M_\chi^2)^2} \right\}.$$

The fermionic dark matter ψ can annihilate into the scalar dark matter H_0 when kinematically favoured. Such conversion are relevant in situations where H_0 is underabundant, as in Scenario 2 discussed in Section 2.2.2.

$$\sigma_{\psi\bar{\psi} \rightarrow H_0 H_0} = \frac{y_3^2}{128\pi} \frac{\beta_\psi \beta_{H_0}}{s} \left\{ \frac{\lambda_H^2 \sin^2 \alpha}{(1 - \tau_H)^2 + \omega_H^2} + \frac{\lambda_{H_S}^2 \cos^2 \alpha}{(1 - \tau_{H_S})^2 + \omega_{H_S}^2} + \lambda_{H_S} \lambda_H \sin 2\alpha \frac{(1 - \tau_{H_S})(1 - \tau_H) + \omega_H \omega_{H_S}}{[(1 - \tau_H)^2 + \omega_H^2][(1 - \tau_{H_S})^2 + \omega_{H_S}^2]} \right\} \quad (2.A.2)$$

where $\lambda_H = (\mu_5/\sqrt{2}) \sin \alpha - 2\lambda_L v \cos \alpha$, $\lambda_{H_S} = (\mu_5/\sqrt{2}) \cos \alpha + 2\lambda_L v \sin \alpha$, $\tau_i = M_i^2/s$, $\omega_i = \Gamma_i M_i/s$ and $\beta_i = \sqrt{1 - 4M_i^2/s}$.

The cross sections for ψ pair annihilation into SM particles relevant to discussion in Section 2.2.2 are given below.

$$\sigma_{\psi\bar{\psi} \rightarrow WW} = \frac{y_3^2}{64\pi} \frac{m_W^4}{v^2} \frac{\beta_\psi \beta_W}{s} \left[3 - \frac{s}{M_W^2} + \frac{s^2}{4M_W^4} \right] \sin^2 2\alpha \left\{ \frac{(\tau_H - \tau_{H_S})^2 + (\omega_H - \omega_{H_S})^2}{[(1 - \tau_H)^2 + \omega_H^2][(1 - \tau_{H_S})^2 + \omega_{H_S}^2]} \right\}, \quad (2.A.3)$$

$$\sigma_{\psi\bar{\psi} \rightarrow ZZ} = \frac{y_3^2}{128\pi} \frac{m_Z^4}{v^2} \frac{\beta_\psi \beta_Z}{s} \left[3 - \frac{s}{M_Z^2} + \frac{s^2}{4M_Z^4} \right] \sin^2 2\alpha \left\{ \frac{(\tau_H - \tau_{H_S})^2 + (\omega_H - \omega_{H_S})^2}{[(1 - \tau_H)^2 + \omega_H^2][(1 - \tau_{H_S})^2 + \omega_{H_S}^2]} \right\}, \quad (2.A.4)$$

$$\sigma_{\psi\bar{\psi} \rightarrow f\bar{f}} = N_C \frac{y_3^2}{8\pi} \frac{m_f^2}{v^2} \frac{\beta_\psi \beta_f^3}{s} \sin^2 2\alpha \left\{ \frac{(\tau_H - \tau_{H_S})^2 + (\omega_H - \omega_{H_S})^2}{[(1 - \tau_H)^2 + \omega_H^2][(1 - \tau_{H_S})^2 + \omega_{H_S}^2]} \right\}, \quad (2.A.5)$$

where $N_C = 1$ for leptons and 3 for quarks.

Cross section for ψ pair annihilating into H_S pair is given by

$$\sigma_{\psi\bar{\psi} \rightarrow H_S H_S} = \frac{1}{16\pi s^2} \frac{1}{\beta_\psi^2} \int_{t_{min}}^{t_{max}} |\mathcal{M}_{H_S}^2| dt \quad (2.A.6)$$

where $t_{min} = -(s/4)(\beta_{H_S} + \beta_\psi)^2$, $t_{max} = -(s/4)(\beta_{H_S} - \beta_\psi)^2$ and the invariant amplitude square is

$$|\mathcal{M}_{H_S}^2| = \frac{y_3^4 \cos^4 \alpha}{8} \left[\frac{M_{H_S}^4 + 2M_{H_S}^2(M_\psi^2 - t) - (4M_\psi^2 - s - t)(3M_\psi^2 + t)}{(t - M_\psi^2)(u - M_\psi^2)} - \frac{t^2 + (s - 2M_{H_S}^2)t - M_\psi^2(s - 7t + 6M_{H_S}^2) + M_{H_S}^4 + 8M_\psi^4}{(t - M_\psi^2)^2} \right] \quad (2.A.7)$$

Cross section for ψ annihilating into a pair of χ is

$$\sigma_{\psi\bar{\psi} \rightarrow \chi^+ \chi^-} = \frac{y_3^2 y_2^2}{64\pi s} \beta_\psi \beta_\chi^3 \left\{ \frac{\sin^4 \alpha}{(1 - \tau_H)^2 + \omega_H^2} + \frac{\cos^4 \alpha}{(1 - \tau_{H_S})^2 + \omega_{H_S}^2} + 2 \sin^2 \alpha \cos^2 \alpha \frac{\omega_H \omega_{H_S} + (1 - \tau_H)^2 (1 - \tau_{H_S})^2}{[(1 - \tau_H)^2 + \omega_H^2][(1 - \tau_{H_S})^2 + \omega_{H_S}^2]} \right\} \quad (2.A.8)$$

Probing multicomponent extension of inert doublet model with a vector dark matter

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3.1 Motivation

In this chapter, we consider another vector boson DM candidate in addition to the existing IDM. While the IDM is odd under a discrete $\mathbb{Z}_2$ symmetry, the added vector boson is also considered to be odd under another $\mathbb{Z}'_2$ symmetry making both the candidates stable. There are other possible ways to stabilise the vector such as the inclusion of a dark $U(1)_D$ symmetry [17, 181]. Such scenarios include an extra scalar field that mixes with the SM Higgs boson after symmetry breaking and provides a stable vector boson dark matter. Here, we consider a minimal scenario avoiding such mixings and stable vector DM is achieved only by introducing a $\mathbb{Z}_2$ symmetry. We have observed that in certain conditions it is even possible to boost the relic density contribution from inert doublet dark matter. This multicomponent scenario, therefore, provides a window to explore the inert dark matter in the mass range 130-500 GeV. Besides, the inclusion of a vector dark matter candidate enriches the dark sector. In this work, we present the phenomenology of the proposed multi-component dark matter model and test the viability of the model by constraining the model parameter space from different theoretical and experimental bounds. We investigate the possibility whether the relic density of inert dark matter candidate can be enhanced significantly in the intermediate regime of inert doublet mass ~ 130 -500 GeV. The proposed multi-component dark matter model further opens up a new window of IDM to be tested in the LHC for possible signatures and promising outcomes.

3.2 Model

In the present work, we consider a multi-component dark matter model by including an additional vector boson to the Inert Doublet Model (IDM) of dark matter. The inert doublet is considered to be odd under a $\mathbb{Z}_2$ which ensures that it does not acquire any vacuum expectation value (v_{ev}) after spontaneous symmetry breaking. Also, particles in the Standard Model (SM) are even under this $\mathbb{Z}_2$ which forbids the decay of the lightest inert particle. Therefore, the lightest inert particle can serve as a dark matter (DM) candidate. Similarly, we impose another discrete $\mathbb{Z}'_2$ symmetry upon the vector boson X which can also be a feasible DM candidate in our model. We denote the SM Higgs doublet as H while the inert Higgs doublet is Φ . The total Lagrangian of our study is as follows

$$\mathcal{L} = |D_\mu H|^2 + |D_\mu \Phi|^2 - V_{\text{IDM}} + \mathcal{L}_{H/\Phi, X} + \mathcal{L}_X + \mathcal{L}_{\text{Yuk}} + \mathcal{L}_{\text{SM}}^{\text{Higgs}}, \quad (3.2.1)$$

where $\mathcal{L}_{\text{SM}}^{\text{Higgs}}$ is the SM Lagrangian sans the Higgs part, $\mathcal{L}_{\text{Yuk}}$ is the SM Yukawa interaction Lagrangian. $\mathcal{L}_X$ is the vector DM sector which is given by

$$\mathcal{L}_X = -\frac{1}{4}X^{\mu\nu}X_{\mu\nu} + \frac{1}{2}M_X^2 X_\mu X^\mu + \frac{1}{4}\lambda_X(X_\mu X^\mu)^2, \quad (3.2.2)$$

the interaction term, $\mathcal{L}_{H/\Phi, X}$ between the scalar fields H, Φ and the dark vector boson, X is¹

$$\mathcal{L}_{H/\Phi, X} = \lambda_{HX}(X^\mu X_\mu)(H^\dagger H) + \lambda_{\Phi X}(X^\mu X_\mu)(\Phi^\dagger \Phi); \quad (3.2.3)$$

and finally, the well-known IDM potential

$$V_{\text{IDM}} = m_{11}^2(H^\dagger H) + m_{22}^2(\Phi^\dagger \Phi) + \lambda_1(H^\dagger H)^2 + \lambda_2(\Phi^\dagger \Phi)^2 + \lambda_3(H^\dagger H)(\Phi^\dagger \Phi) + \lambda_4(H^\dagger \Phi)(\Phi^\dagger H) + \frac{\lambda_5}{2}[(H^\dagger \Phi)^2 + \text{h.c.}], \quad (3.2.4)$$

where all the couplings $\lambda_i, i = 1 - 5$ are assumed to be real and also $m_{11}^2 < 0$ so that spontaneous symmetry breaking (SSB) occurs. After SSB the SM Higgs doublet receives a v_{ev} $v = 246$ GeV and the doublet fields are denoted as

$$H = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}(v + h) \end{pmatrix}, \quad \Phi = \begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}(H_0 + iA_0) \end{pmatrix}. \quad (3.2.5)$$

Masses of different physical scalar including SM Higgs and inert particles and the vector boson X of the dark sector are given as follows

$$\begin{aligned} m_h^2 &= 2\lambda_1 v^2 \\ m_{H^\pm}^2 &= m_{22}^2 + \lambda_3 \frac{v^2}{2} \\ m_{H_0}^2 &= m_{22}^2 + (\lambda_3 + \lambda_4 + \lambda_5) \frac{v^2}{2} \\ m_{A_0}^2 &= m_{22}^2 + (\lambda_3 + \lambda_4 - \lambda_5) \frac{v^2}{2} \\ m_X^2 &= M_X^2 + \lambda_{HX} v^2. \end{aligned} \quad (3.2.6)$$

¹ $(X^\mu X_\mu)(H^\dagger H)$ and $H^2 \Phi^2$ terms may contribute to the Higgs effective potential. The new physics contribution to the Higgs thermal mass is dependent on $\lambda_L, \lambda_S (= (\lambda_3 + \lambda_4 - \lambda_5)/2)$ and λ_3 in the IDM sector and λ_{HX} in the X sector. With these parameters, it may be possible to obtain a strongly first order phase transition, as shown in [74, 95]. However, this is not relevant to our discussion. Without affecting the phenomenology, one can set $\lambda_{HX} = 0$ and therefore eliminate the effect of dark vector boson in the Higgs effective potential.

In the above Eq. (3.2.6), m_h refers to mass of SM Higgs boson (h) $m_h = 125.09$ GeV [225]. As mentioned, even after SSB the dark $\mathbb{Z}_2$ ($\mathbb{Z}'_2$) symmetry of inert doublet (vector boson X) remains intact and the lightest scalar H_0 of the IDM² along with X provides a scenario for multicomponent dark matter.

Before we present the discussions on the phenomenology of multicomponent dark matter model, we first mention some important theoretical and experimental bounds to be taken into account in the next section.

3.3 Constraints and bounds

Vacuum Stability: In order to stabilise the Higgs vacuum following conditions must be satisfied [176]

$$\begin{aligned} \lambda_1, \lambda_2 > 0, \lambda_3 + 2\sqrt{\lambda_1\lambda_2} > 0, \\ \lambda_3 + \lambda_4 - |\lambda_5| + 2\sqrt{\lambda_1\lambda_2} > 0. \end{aligned} \quad (3.3.1)$$

Perturbativity: Quartic interactions at tree level at high energy provides bound on the eigenvalues $|\Lambda_i|$ of quartic couplings which must obey the relation $|\Lambda_i| \leq 4\pi$.

LEP: LEP [225] provides bound from the decay width measurement of Z boson which is given as

$$\begin{aligned} m_{H_0} + m_{A_0} > m_Z, \\ m_{H^\pm} > 80 \text{ GeV}. \end{aligned} \quad (3.3.2)$$

Relic Density: In the present model we have two feasible dark matter candidates H_0 and X . Since both the candidates contribute to the dark matter relic density obtained from Planck [16] experiment, one must satisfy the following relation

$$\Omega_{\text{DM}} h^2 = 0.1199 \pm 0.0027 = \Omega_{H_0} h^2 + \Omega_X h^2, \quad (3.3.3)$$

where h denotes the Hubble parameter ($100 \text{ km s}^{-1} \text{ Mpc}^{-1}$) and relic density of inert doublet (vector dark matter) is given as $\Omega_{H_0} h^2$ ($\Omega_X h^2$).

Direct Detection Experiments: Apart from satisfying the conditions from dark matter relic density, both the dark matter candidates must agree with present direct detection bounds from different dark matter search experiments as well. In this work, we constrain the model parameter space with the latest direct detection limits obtained from XENON1T [29].

Indirect Detection Experiments: DM is further constrained from the observations of diffused γ -rays from the Galactic Centre (GC) and dwarf spheroidal galaxies (dSphs) where DM density appears to be large. Experiments such as Fermi-LAT [20] and H.E.S.S [13] has put constraints on the upper limit of velocity averaged scattering cross sections for various channels which can contribute to the observed photon flux. Here we obeyed constraints on the model parameter space emanating from the observations of the latest Fermi-LAT data [20].

Searches at LHC: The DM searches at the LHC has been performed in various hadronic as well as leptonic channels. Here we will follow the *dilepton + missing energy* ($2\ell + \cancel{E}_T$) searches for our present study. Usually, such experimental searches were conducted in the context of Supersymmetric (SUSY) theories. Since our study does not include SUSY, we will follow the outcome of those searches with some care.

3.4 Formalism for Dark Matter Analysis

Before we begin our analysis of the multi-component dark matter scenario, we briefly mention the calculations of relic density and direct detection measurements in the present model.

²We consider $\lambda_5 < 0$.

3.4.1 Relic density of dark matter candidates

Since the present model deals with two dark matter candidates which also interact with themselves, we have to solve for the coupled Boltzmann equation. Relic density for each of the dark matter candidate is obtained by solving these coupled equations which are written as

$$\frac{dn_{H_0}}{dt} + 3Hn_{H_0} = -\langle\sigma v\rangle_{H_0H_0 \rightarrow SM SM} (n_{H_0}^2 - n_{H_0eq}^2) + \langle\sigma v\rangle_{XX \rightarrow H_0H_0} \left(n_X^2 - \frac{n_{Xeq}^2}{n_{H_0eq}^2} n_{H_0}^2 \right); \quad (3.4.1)$$

$$\frac{dn_X}{dt} + 3Hn_X = -\langle\sigma v\rangle_{XX \rightarrow SM SM} (n_X^2 - n_{Xeq}^2) - \langle\sigma v\rangle_{XX \rightarrow H_0H_0} \left(n_X^2 - \frac{n_{Xeq}^2}{n_{H_0eq}^2} n_{H_0}^2 \right). \quad (3.4.2)$$

where $m_X > m_{H_0}$. Similarly, for $m_{H_0} > m_X$,

$$\frac{dn_{H_0}}{dt} + 3Hn_{H_0} = -\langle\sigma v\rangle_{H_0H_0 \rightarrow SM SM} (n_{H_0}^2 - n_{H_0eq}^2) - \langle\sigma v\rangle_{H_0H_0 \rightarrow XX} \left(n_{H_0}^2 - \frac{n_{H_0eq}^2}{n_{Xeq}^2} n_X^2 \right); \quad (3.4.3)$$

$$\frac{dn_X}{dt} + 3Hn_X = -\langle\sigma v\rangle_{XX \rightarrow SM SM} (n_X^2 - n_{Xeq}^2) + \langle\sigma v\rangle_{H_0H_0 \rightarrow XX} \left(n_{H_0}^2 - \frac{n_{H_0eq}^2}{n_{Xeq}^2} n_X^2 \right). \quad (3.4.4)$$

In the above Eqs. (3.4.1) and (3.4.2), $n_i, i = X, H_0$, denotes the number density of dark matter particles and their equilibrium number densities are expressed as n_{ieq} respectively. Annihilation cross section of dark matter candidates into SM are given as $\langle\sigma v\rangle_{ii \rightarrow SM SM}$ while the same between themselves are denoted by the $\langle\sigma v\rangle_{XX \rightarrow H_0H_0}$ ($\langle\sigma v\rangle_{H_0H_0 \rightarrow XX}$) for $m_X > m_{H_0}$ ($m_{H_0} > m_X$). Solving for the Eqs. (3.4.1) and (3.4.2) (or Eqs. (3.4.3) and (3.4.4) depending on the masses of m_{H_0} and m_X), one can obtain the relic density contributions from each of the dark matter candidates of the form

$$\Omega_i h^2 = 2.755 \times 10^8 \frac{m_i}{GeV} Y_i(T_0), \quad i = H_0, X, \quad (3.4.5)$$

where $Y_i = n_i/s$, is the yield of dark matter candidate obtained at a present temperature of Universe T_0 and s is the entropy density of the Universe. The total DM relic density is then obtained by adding individual relic densities of both the candidates as mentioned in Eq. (3.3.3). It is to be noted that the annihilation cross section $\langle\sigma v\rangle_{XX \rightarrow H_0H_0}$ ($\langle\sigma v\rangle_{H_0H_0 \rightarrow XX}$) depends on the coupling between dark sector particles and therefore the coupling $\lambda_{\Phi X}$ plays a significant role in the dark matter phenomenology.

3.4.2 Direct detection of DM candidates

Dark matter direct search experiments like LUX, XENON1T etc. search for direct interactions of dark matter with detector nuclei. Dark matter candidate can undergo elastic scattering with detector nuclei and recoil energy will be transferred which can be measured at the detector. However, no such event has been recorded yet which provides a stringent limit on dark matter direct detection cross section. In the present model, both of dark matter candidates H_0 and X can undergo spin-independent elastic scattering with the detector nuclei. Since, the model involves two dark matter candidates, the final direct detection cross section for each dark matter will be scaled by a factor $r_i = \Omega_i h^2 / \Omega_{DM} h^2, i = H_0, X$. Bounds from direct detection of dark matter will constrain the model parameters. It is to be noted that coupling $\lambda_{\Phi X}$ has no contribution to the direct detection measurements. The spin-independent direct detection cross section for the scalar dark matter H_0 is given as

$$\sigma_{H_0}^{SI} = r_{H_0} \frac{\lambda_L^2}{16\pi} \frac{1}{m_h^4} f^2 \frac{m_N^4}{(m_{H_0} + m_N)^2}, \quad (3.4.6)$$

where $\lambda_L = (\lambda_3 + \lambda_4 + \lambda_5)/2$ and m_N denotes the mass of the nucleon. In the above expression of Eq. (3.4.6) the contributions from nuclear matrix elements are given by the factor $f \sim 0.3$ [19, 47]. Similarly, the spin-independent direct detection cross section for the vector dark matter candidate X is expressed as

$$\sigma_X^{\text{SI}} = r_X \frac{\lambda_{HX}^2}{16\pi} \frac{1}{m_h^4} f^2 \frac{m_N^4}{(m_X + m_N)^2}, \quad (3.4.7)$$

We constrain the model parameter space using the most stringent direct detection bounds obtained from XENON1T [29].

3.5 Observations and results

Before we present the discussions on our model with multi-component dark matter, we summarise the main parameters in the model. The model is implemented in FeynRules [23] and the relic density computation is performed using micrOMEGAS [50] package scanning over the available parameter space. As mentioned previously, the model has two dark matter candidates, the lightest inert doublet particle H_0 and the vector boson X . The parameters from the inert doublet that contribute to relic density and direct detection measurements are well known

$$\lambda_L, m_{H_0}, m_{A_0}, m_{H^\pm}.$$

Similarly, the parameters that contribute to DM phenomenology of the vector boson are

$$\lambda_{HX}, m_X.$$

Apart from these above-mentioned parameters, there is another coupling $\lambda_{\Phi X}$ which can contribute to the annihilation among the dark sector particles. It is to be noted that even for the coupling $\lambda_{\Phi X} = 0$, the above annihilation can occur through Higgs mediated diagrams. However, such contributions are significant only near the Higgs resonance and become small when we consider the mass of dark matter candidates away from Higgs resonance. Also, the strong bound from dark matter direct detection severely constrains the couplings λ_L and λ_{HX} reducing the effects of dark sector annihilation. On the other hand, the four-point coupling $\lambda_{\Phi X}$ is a completely independent parameter which does not contribute to dark matter direct detection but can contribute to relic density of DM particles. In this work, our primary aim is to study how the $2 \leftrightarrow 2$ annihilation between dark sector particles H_0 and X affects the multi-component dark matter scenario. To this end, we consider the intermediate regimes of dark matter masses in our work

$$100 \text{ GeV} \leq m_{H_0}, m_X \leq 500 \text{ GeV}. \quad (3.5.1)$$

It is to be noted that for the case of the inert doublet, co-annihilation effects can be significant if mass splitting between the scalar H_0 , A_0 and the charged particle $H^\pm$ is small. In fact, for a pure inert doublet dark matter, relic density is negligible for large mass splitting. In this work, we present our results for two values of mass splitting $\Delta m = 10 \text{ GeV}$ and 25 GeV where $\Delta m = m_{A_0} - m_{H_0}$ (in GeV) and $m_{H^\pm} = m_{A_0} + 0.1 \text{ GeV}$ ³. The coupling λ_L and λ_{HX} should not be large in order to satisfy direct detection bounds and we set them to be equal, $\lambda_L = \lambda_{HX} = 0.01$. In this way, we restrict other model parameters and use the coupling $\lambda_{\Phi X}$ as a variable, the new parameter to determine the allowed regions in this framework.

To begin with, we first consider a simplified case setting $\lambda_L = \lambda_{HX} = 0$ and changing $\lambda_{\Phi X}$ from a very small value to 0.1 ⁴. We denote the relic density of inert dark matter H_0 as $\Omega_{H_0} h^2$ and that of the vector DM candidate as $\Omega_X h^2$ (as expressed in Eq. (3.3.3)).

³ $m_{H^\pm} - m_{A_0} = 0.1 \text{ GeV}$ is maintained throughout the analysis for all the different cases considered in the work.

⁴It is to be noted that since we have already set the coupling $\lambda_{HX} = 0$, we cannot use $\lambda_{\Phi X} = 0$ when solving the coupled Boltzmann equation. Hence we considered a very small $\lambda_{\Phi X} \sim 10^{-15}$ such that it reproduces the nature of IDM.

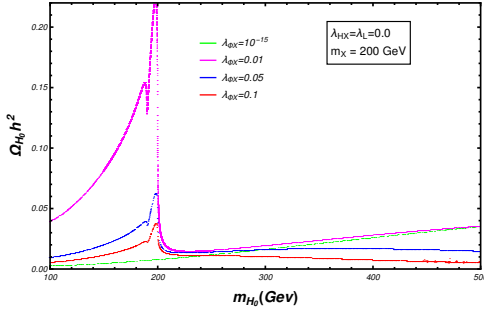


Figure 3.5.1: $\Omega_{H_0} h^2$ vs m_{H_0} for $\Delta m = 10$ GeV. Only the quartic interactions in $XX \rightarrow H_0 H_0$ and quartic gauge interactions in $H_0 H_0 \rightarrow SM SM$ are considered.

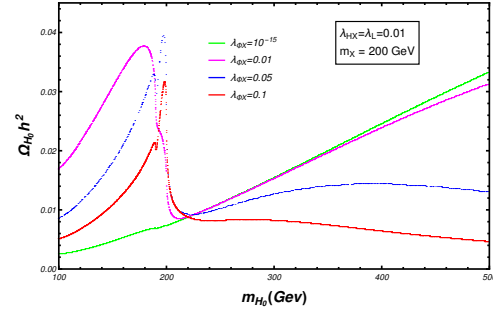


Figure 3.5.2: $\Omega_{H_0} h^2$ vs m_{H_0} for $\Delta m = 10$ GeV. All the possible interactions (both quartic and Higgs mediated) in $H_0 H_0 \rightarrow SM SM$ and $XX \rightarrow H_0 H_0$ are taken into account.

In Fig. 3.5.1 we show the variation of $\Omega_{H_0} h^2$ against m_{H_0} for the chosen values of couplings $\lambda_{\Phi X}$ for $\lambda_L = \lambda_{HX} = 0$. With this consideration, only the gauge interactions of inert dark matter will survive and there will be $H_0 H_0 \leftrightarrow XX$ annihilation due to quartic coupling $\lambda_{\Phi X}$ only. Moreover, with $\lambda_L = \lambda_{HX} = 0$, there will not be any direct detection signatures of both the dark matter candidates. We set the mass difference $\Delta m = m_A - m_{H_0} = 10$ GeV and $m_X = 200$ GeV. From Fig. 3.5.1, we observe that for $\lambda_{\Phi X} = 0.01$, the relic density $\Omega_{H_0} h^2$ of H_0 changes drastically from the case with $\lambda_{\Phi X} = 10^{-15}$, where H_0 can annihilate only to SM gauge bosons for the region of mass $m_{H_0} \leq 200$ GeV and afterwards follows the same pattern of IDM ($\lambda_{\Phi X} = 10^{-15}$). The increase in H_0 relic density happens till $m_{H_0} \sim m_X$ and it falls sharply as $m_{H_0} \geq m_X$ when $H_0 H_0 \rightarrow XX$ annihilation channel opens. Apart from that there exists another small peak around $m_{H_0} \sim 190$ GeV as new annihilation channels $A_0 A_0, H^+ H^- \rightarrow XX$ appear before $H_0 H_0 \rightarrow XX$ annihilation. Comparing the plots in Fig. 3.5.1 for the case $\lambda_{\Phi X} = 10^{-15}$ with $\lambda_{\Phi X} = 0.01$, we conclude that the enhancement in the relic density is due to the production of H_0 particles from annihilations $XX \rightarrow H_0 H_0$. However, with increasing $\lambda_{\Phi X}$, this situation changes. With increase in $\lambda_{\Phi X}$, the annihilation $\langle \sigma v \rangle_{XX \rightarrow H_0 H_0}$ also increases. However, this results in a reduction of n_X number density and rise in n_{H_0} number density. As a result, the overall contribution of the second term in Eq. (3.4.1) decreases which is clearly observed in Fig. 3.5.1. Therefore, relic density of H_0 (as well as enhancement in relic density) decreases with increasing $\lambda_{\Phi X}$ coupling in the region $m_{H_0} < m_X \sim 200$ GeV and tends to decrease further in the higher mass range $m_{H_0} > m_X$ governed by Eq. (3.4.3).

A discussion is in order regarding whether inert doublet itself can obtain total DM relic abundance assisted by the vector boson. Looking into Fig. 3.5.1 one may think that it is possible as for $\lambda_{\Phi X} = 0.01$, relic density of IDM is overabundant. However, it is to be noted that this is the case for $\lambda_L = 0$. An interesting feature of IDM that distinguishes it from ordinary scalar DM is the very strong gauge interactions. As a result, even with $\lambda_L = 0$ and no conversion into vector particle via direct interaction ($\lambda_{\Phi X} = 0$) and Higgs portal (since $\lambda_{HX}, \lambda_L = 0$), for $m_{H_0} \geq m_{W,Z}$ we have large annihilation into gauge bosons and DM relic density is very low following the green plot shown in Fig. 3.5.1. Now if we consider a pure IDM case with $\lambda_L \neq 0, \lambda_{HX} = \lambda_{\Phi X} = 0$, with new annihilation channels into SM particles relic density of IDM will decrease further. However, in the Higgs portal DM case, since there is no direct interaction with the gauge sector, the annihilation of DM depends only on Higgs portal coupling and with an increase in the coupling, relic density decreases. Since in IDM, the gauge interaction dominates for $m_{H_0} > m_{W,Z}$ a small value of $\lambda_L = 0.01$ will not affect the relic abundance very much but for a higher value of $\lambda_L = 0.1$, the relic density can be even smaller due to more annihilations. Also, larger λ_L values will eliminate most of the parameter space (for low mass IDM $m_{H_0} < 500$ GeV) due to large direct detection cross section. However, in the presence of the vector DM candidate, it is possible to enhance the IDM relic density from conversion mechanism with new production channels $XX \rightarrow H_0 H_0$ as discussed in Fig. 3.5.1. Now, we will discuss a case of general IDM with $\lambda_L \neq 0$ and try to investigate to what extent the relic density of IDM can be enhanced in the

general situation.

We now consider a more realistic picture than what is shown in Fig. 3.5.1, with all possible channels of inert dark matter annihilation into consideration by taking $\lambda_L \neq 0$. For this, we demonstrate our results with fixed Yukawa couplings $\lambda_L, \lambda_{HX} = 0.01$ with $\Delta m = 10$ GeV and $m_X = 200$ GeV for four values of $\lambda_{\Phi X} = 10^{-15}, 0.01, 0.05, 0.1$.⁵ We show the variation in relic density of inert scalar dark matter H_0 with mass m_{H_0} in Fig. 3.5.2. From Fig. 3.5.2 we observe that for $\lambda_{\Phi X} = 10^{-15}$, the results for inert scalar dark matter are identical to the usual inert dark matter model. However, situations change as the mixing $\lambda_{\Phi X}$ becomes large and a prominent resonance appears near the mass of m_X and relic density enhances as m_{H_0} approaches m_X . Plots with different $\lambda_{\Phi X} = 0.01, 0.05, 0.1$ follow the similar pattern in appearing in Fig. 3.5.1. However, due to the presence of new annihilation channels of H_0 (as $\lambda_L = 0.01$), the enhancement in $\Omega_{H_0} h^2$ is reduced for the mass range $m_{H_0} \leq m_X$ when compared with Fig. 3.5.1. For the regime, $m_{H_0} \leq m_X$, inert scalar particles are being produced via the annihilation of X particles which compensates the annihilation of $H_0 H_0 \rightarrow W^+ W^-$. Therefore although H_0 particles are being annihilated they are also produced at a larger rate which enhances their relic density contribution from IDM. It is to be noted that although there is a new production channel $XX \rightarrow H_0 H_0$ through Higgs, its contribution is not significant and suppressed due to direct production depending on coupling $\lambda_{\Phi X}$. The inert scalar relic density then decreases with increasing mass $m_{H_0} > m_X$ and this effect becomes large for larger values of $\lambda_{\Phi X}$. This indicates that for the region of mass $m_{H_0} > m_X$, $H_0 H_0 \rightarrow XX$ annihilation becomes large and as a result relic density of IDM candidate reduces considerably. The most interesting feature is the resonance region where IDM relic density is enhanced within the 100-500 GeV regime depending on the mass of m_X (which is required to be in the intermediate-mass regime as well) when compared with the normal single component IDM scenario equivalent to the case $\lambda_{\Phi X} = 10^{-15}$. As seen in Fig. 3.5.2, relic density of inert dark matter candidate H_0 is $\Omega_{H_0} h^2 \sim 0.032$ for $m_{H_0} = 200$ GeV with $\lambda_{\Phi X} = 0.01$ which is much larger than $\Omega_{H_0} h^2 = 7.8 \times 10^{-3}$ when compared with the usual single component IDM case (for $\lambda_{\Phi X} = 10^{-15}$ shown in Fig. 3.5.1). Moreover, the same order of relic density is achieved at mass $m_{H_0} \sim 500$ GeV in the normal inert doublet. We also observe that increasing $\lambda_{\Phi X}$ also reduces DM relic density $\Omega_{H_0} h^2$ by increasing contribution of $H_0 H_0 \leftrightarrow XX$ channel which then starts to dominate over inert doublet annihilations in Boltzmann equation following the same pattern obtained in Fig. 3.5.1. Therefore, one can have a significant contribution from inert DM candidate in the present scenario even within the mass range 130-500 GeV.

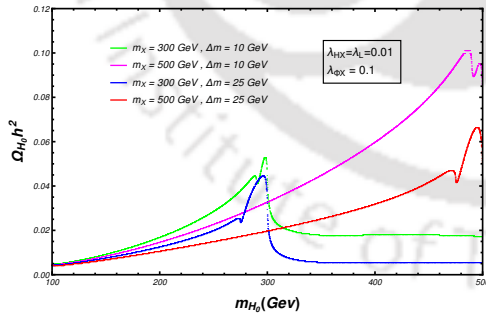


Figure 3.5.3: $\Omega_{H_0} h^2$ vs m_{H_0} for $\lambda_{\Phi X} = 0.1$ as a function of m_X .

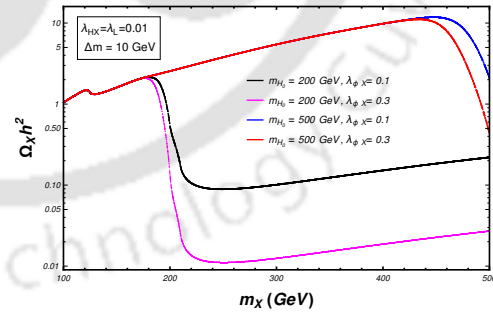


Figure 3.5.4: $\Omega_X h^2$ vs m_X for different values of m_{H_0} and $\lambda_{\Phi X}$.

In another plot (Fig. 3.5.3) we show the variations of $\Omega_{H_0} h^2$ for the same range of m_{H_0} with two different values of $m_X = 300$ GeV and 500 GeV for $\lambda_{\Phi X} = 0.1$ taking $\Delta m = 10$ GeV and compare the results with $\Delta m = 25$ GeV for same set of parameters. The couplings λ_L and λ_{HX} are kept fixed at the same value 0.01. We observe that for $m_X = 300$ GeV, $\Omega_{H_0} h^2$ increases with m_{H_0} until the resonance. In this regime, the annihilation $XX \rightarrow H_0 H_0$ dominates which enhances $\Omega_{H_0} h^2$ and a

⁵As mentioned earlier, for larger values of $\lambda_L = 0.1$, there will be large annihilations of IDM which will further decrease the IDM relic density. Moreover, larger λ_L will also result in large direct detection cross section and reduce low mass IDM parameter space. Therefore we consider a smaller value of $\lambda_L = 0.01$. A detailed discussion of λ_L is presented later.

sharp fall appears thereafter. Similar behaviour is observed for $m_X = 500$ GeV. This nature is followed for both values of Δm . However, it is to be noted that for the case $\Delta m = 10$ GeV, $\Omega_{H_0} h^2$ is larger than when $\Delta m = 25$ GeV. This is because, for smaller Δm , the contribution of co-annihilation channels in inert doublet is large which can significantly enhance relic density of inert scalar H_0 . On the other hand for larger mass splitting, co-annihilation channels are suppressed resulting reduction of relic abundance. The spread of the peak near $m_{H_0} \sim m_X$ also depends on the value of mass splitting as clearly observed in Fig. 3.5.3. The resonances peak in the plots is sharper for $\Delta m = 10$ GeV and broader for $\Delta m = 25$ GeV. The resonance peak near $m_{H_0} \sim m_X$ is associated with another small peak which appears as $A_0 A_0 \leftrightarrow XX$ and $H^\pm H^\pm \leftrightarrow XX$ channels open up before $H_0 H_0 \leftrightarrow XX$. Therefore, one can conclude that depending on the mass of vector DM candidate and mass splitting Δm , H_0 can have a formidable contribution to total DM relic density. However, from both Figs. 3.5.2 and 3, we observe that although relic abundance of inert dark matter is increased, it cannot alone satisfy the total DM relic density for dark matter and therefore the vector dark matter also contributes to the total DM relic density.

So far we have discussed how relic abundance of IDM changes with the inclusion of additional vector dark matter and found that although there is a sizeable enhancement in H_0 relic density, we need contribution from vector DM also. It is to be noted that vector DM can alone satisfy the DM relic density (see Ref. [39] and references therein). However, it is found that although portal vector DM can explain DM relic density with $\lambda_{HX} \sim 0.1$, it fails to satisfy direct detection limits for a large range of mass range up to 1 TeV (see Fig. 19 of Ref. [39]). Therefore, we investigate in this work whether a low mass vector DM can be achieved to satisfy direct detection limits and also provide the required amount of relic abundance such that Eq. (3.3.3) is satisfied. If the coupling λ_{HX} is increased further, the relic abundance of pure Higgs portal vector dark matter will fall (as $\Omega h^2 \sim 1/\langle\sigma v\rangle$), but it will be ruled out by direct detection. On the other hand for smaller $\lambda_{HX} = 0.01$, direct detection can be recovered but vector DM becomes overabundant. Hence, it is interesting to study whether, in the two-component framework, relic density of vector dark matter can be reduced for small λ_{HX} .

In Fig. 3.5.4, we present the variation of relic density of vector dark matter with its mass m_X for two different values of $\lambda_{\Phi X} = 0.1, 0.3$ and for fixed $\lambda_L = \lambda_{HX} = 0.1$ and $\Delta m = 10$ GeV. With this choice, we can directly investigate whether we have a region where vector dark matter becomes underabundant to satisfy the requirements of combined relic density following Eq. (3.3.3). From Fig. 3.5.4, we observe that for $m_{H_0} = 200$ GeV, the relic abundance of vector DM initially starts with a higher value for $m_X \leq m_{H_0}$. An initial drop in relic density occurs when the annihilation channel $XX \rightarrow hh$ opens for $m_X = m_h$. However, as m_X increases new annihilation channel $XX \rightarrow H_0 H_0$ (when $m_X \sim m_{H_0}$) opens up and as a result relic density of X decreases. This reduction in relic density depends on the coupling $\lambda_{\Phi X}$. It can be easily observed from Fig. 3.5.4 that for $\lambda_{\Phi X} = 0.3$, annihilation $XX \rightarrow H_0 H_0$ is large and relic density of vector DM falls considerably when compared with the case of $\lambda_{\Phi X} = 0.1$. Similar nature of the relic density plots is observed for $m_{H_0} = 500$ GeV, but in this case, the relic abundance of vector DM decreases at higher mass $m_X \sim m_{H_0}$ when the annihilation becomes kinematically allowed. Therefore, the low mass region of vector dark matter which was ruled out by direct detection in standard Higgs portal scenario becomes accessible in the two-component framework and also opens the window for the study of multicomponent dark matter.

Now we present our results for the overall range of parameter space described at the beginning of this Section. In Fig. 3.5.5 we depict a scatter plot of m_{H_0} against $\Omega_{H_0} h^2$ for two values of $\lambda_{\Phi X} = 0.1, 0.3$ varying both m_{H_0} and m_X from 100-500 GeV with $\Delta m = 10$ GeV. We have considered the same values of $\lambda_L, \lambda_{HX} = 0.01$ as stated earlier. We observe a large spread in relic density $\Omega_{H_0} h^2$ with m_{H_0} . We also observe in Fig. 3.5.5 that increasing $\lambda_{\Phi X}$ results in decrease in the relic density $\Omega_{H_0} h^2$ which directly follows the nature of plots in Figs. 3.5.1 and 2. For the chosen $\Delta m = 10$ GeV, we observe that the relic density $\Omega_{H_0} h^2$ attains a maximum 0.1 for $m_{H_0} \sim 500$ GeV for $\lambda_{\Phi X} = 0.1$ and for $\lambda_{\Phi X} = 0.3$ it can reach values up to 0.05. This is due to increase in $\lambda_{\Phi X}$ which increases the $2 \leftrightarrow 2$ annihilations between dark sector particles as stated earlier.

In Fig. 3.5.6 we show the variation of m_X with $\Omega_X h^2$ for the same range of parameters. Here we also observe a similar increase in relic density with m_X as obtained for scalar DM candidate. However,

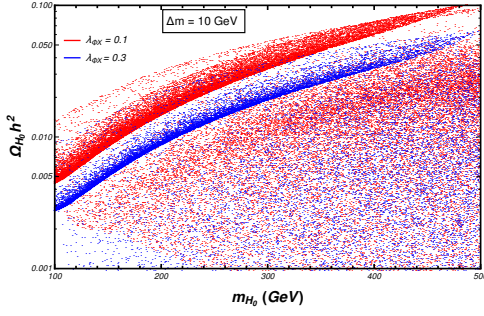


Figure 3.5.5: $\Omega_{H_0} h^2$ vs m_{H_0} as a function of $\lambda_{\Phi X}$ for $\Delta m = 10$ GeV.

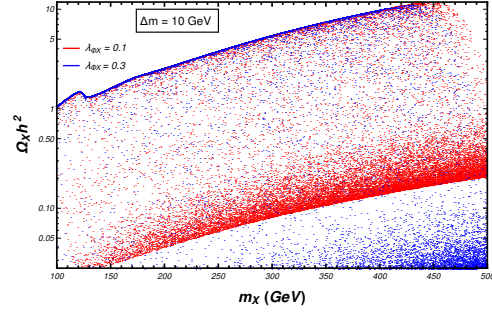


Figure 3.5.6: $\Omega_X h^2$ vs m_X as a function of $\lambda_{\Phi X}$ for $\Delta m = 10$ GeV.

a large range of parameter space is ruled out for as $\Omega_X h^2$ becomes overabundant. A small drop in $\Omega_X h^2$ near $m_X \sim m_h$ represents the new annihilation channels $XX \rightarrow hh$. Comparing Fig. 3.5.4 with Fig. 3.5.6, it can be easily observed that region of upper and lower part of the envelope in Fig. 3.5.6 follows directly from Fig. 3.5.4 as the mass m_{H_0} is varied from 100 to 500 GeV. For smaller values of m_X , as shown in Fig. 3.5.4 a drop in $\Omega_X h^2$ occurs which creates the lower envelope while for higher values of m_X the reduction in relic density appears later creating the upper envelope of the scatter plot. Similar plots for variations for $\Omega_{H_0} h^2$ ($\Omega_X h^2$) against m_{H_0} (m_X) with $\Delta m = 25$ GeV are shown in Figs. 3.5.7 and 8. It is interesting to observe that for increased mass splitting between H_0 and A_0 and for $\lambda_{\Phi X} = 0.1$, the maximum relic density obtained for the inert DM ($\Omega_{H_0} h^2 \sim 0.05$) is almost half when compared with Fig. 3.5.5. This is because as we increase the mass splitting, the contributions from co-annihilation channels will reduce. The plot in Fig. 3.5.7 depicts same nature of Fig. 3.5.5 as we change $\lambda_{\Phi X}$. However, for the vector DM candidate X , there is no such effect and the corresponding Fig. 3.5.8 remains almost similar to Fig. 3.5.6 with no significant change. Both Figs. 3.5.6 and 8 exactly follows the nature of Fig. 3.5.4 discussed earlier. Moreover, the choice of Δm does not affect the relic density of X as it is related to inert dark matter. However, we will show that although it doesn't affect the relic density $\Omega_X h^2$, mass splitting Δm can have a significant effect when the total relic density of the multi-component dark sector is taken into account.

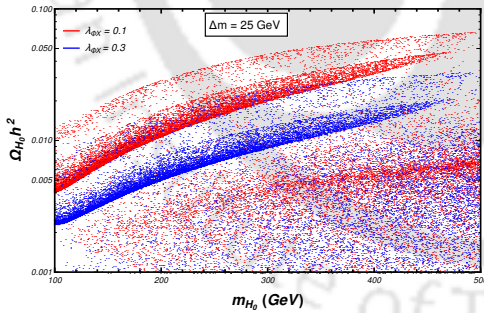


Figure 3.5.7: $\Omega_{H_0} h^2$ vs m_{H_0} as a function of $\lambda_{\Phi X}$ for $\Delta m = 25$ GeV.

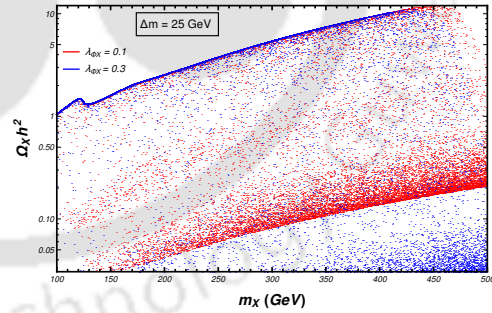


Figure 3.5.8: $\Omega_X h^2$ vs m_X as a function of $\lambda_{\Phi X}$ for $\Delta m = 25$ GeV.

We would like to mention here that a simple scenario for scalar and vector DM without mixing $\lambda_{\Phi X}$ is completely different from inert doublet and vector DM. First of all, inert doublet has direct gauge interaction but singlet scalar DM can interact with gauge boson through Higgs portal only. Also in the case of singlet scalar DM (S), there is only one conversion channel $SS \leftrightarrow XX$ (via quartic coupling $\lambda_{\Phi X}$ and Higgs mediation). But for inert doublet, there are many other annihilation channels, such as $H_0 H_0, A_0 A_0, H^+ H^- \leftrightarrow XX$ (the analytical expressions for the annihilation cross section are given in the Section 3.A), that can affect the result. Furthermore, it also depends on the mass splitting between H_0, A_0 etc. Therefore, results with inert doublet and vector DM are substantially different from the study of singlet scalar + vector DM discussed in Ref. [64]. This can be easily understood from Fig. 3.5.3. In the present multicomponent DM, depending on mass splitting we have two distinct peaks in Fig. 3.5.3,

small peaks due to $A_0 A_0, H^+ H^- \leftrightarrow X X$ and a large peak due to $H_0 H_0 \leftrightarrow X X$. In the case of singlet + vector DM, there will be only one such peak only for $SS \leftrightarrow X X$ as there is no other annihilation possible.

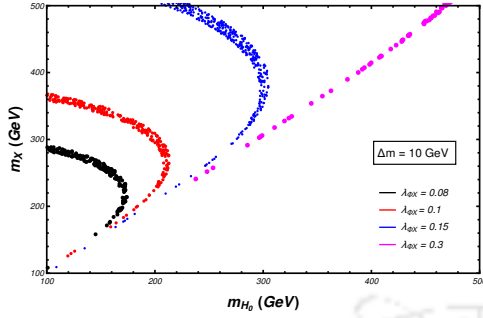


Figure 3.5.9: m_X vs m_{H_0} for $\Delta m = 10$ GeV as a function of $\lambda_{\Phi X}$.

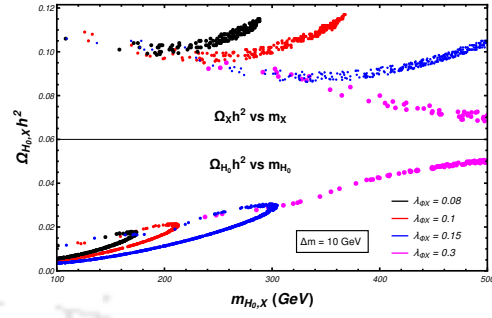


Figure 3.5.10: $\Omega_X h^2$ vs $m_X, \Omega_{H_0} h^2$ vs m_{H_0} for $\Delta m = 10$ GeV.

In Fig. 3.5.9, we present the allowed ranges of dark matter masses in $m_{H_0} - m_X$ plane for four values of $\lambda_{\Phi X} = 0.08, 0.1, 0.15, 0.3$ considering $\lambda_L = \lambda_{HX} = 0.01$ with $\Delta m = 10$ GeV. We vary the masses of dark matter candidates in the range $100 \text{ GeV} \leq m_{H_0, X} \leq 500 \text{ GeV}$ as mentioned in the beginning of this section and use the condition expressed in Eq. (3.3.3) for total DM relic abundance. From the plots in Fig. 3.5.9 we observe that apart from the region when $m_{H_0} \sim m_X$, there exists another regime with $m_X > m_{H_0}$. However this nature disappears as we increase the coupling $\lambda_{\Phi X}$. From the Fig. 3.5.9, we notice that in case of $\lambda_{\Phi X} = 0.1$, the maximum mass $m_{H_0} \sim 220$ GeV and then it tends to decrease with increasing m_X while m_{H_0} reaches a maximum value about 300 GeV for $\lambda_{\Phi X} = 0.15$.

In order to explain the nature of these plots in Fig. 3.5.10 we present the variation of $\Omega_{H_0, X} h^2$ vs $m_{H_0, X}$ for the same set of parameters considered in Fig. 3.5.9. From Fig. 3.5.10, it can be observed that for $\lambda_{\Phi X} = 0.1$, initially the relic density $\Omega_X h^2$ tends to decrease with increasing m_X . In this regime, the contribution of inert dark matter H_0 as well as the mass of H_0 also increases. This can easily be checked by adding up the relic density $\Omega_{H_0} h^2$ with $\Omega_X h^2$ that provide the required total DM relic abundance. This corresponds to the regime $m_X \sim m_{H_0}$, the lower half of the semicircular arc in Fig. 3.5.9 for $\lambda_{\Phi X} = 0.1$ that continues till $m_X \sim m_{H_0} \sim 220$ GeV. For larger values of $m_X > 220$ GeV, the relic density $\Omega_X h^2$ tends to increase while the contribution of H_0 in total DM relic decreases which also results in a reduction of the mass of H_0 . This regime corresponds to the upper part of the semicircular allowed region shown in Fig. 3.5.9. This can also be justified by looking into Fig. 3.5.6 where $\Omega_X h^2$ becomes overabundant near $m_X \sim 380$ GeV for $\lambda_{\Phi X} = 0.1$. This indicates that with the increase in m_X , the contribution of the other candidate must be small as $\Omega_X h^2$ approaches $\Omega_{DM} h^2$. Similar conclusion can be drawn for $\lambda_{\Phi X} = 0.08, 0.15$ from Figs. 3.5.9 and 10. However, with increasing values of $\lambda_{\Phi X} = 0.3$ such feature disappears as observed in Fig. 3.5.9. We have found that with the present two component scenario, the relic density contribution of H_0 can be enhanced significantly. For example, in Fig. 3.5.10, the relic density of H_0 having mass $m_{H_0} = 400$ GeV when calculated for $\lambda_{\Phi X} = 0.3$ is about $\Omega_{H_0} h^2 = 0.042$ with respect to the usual value $\Omega_{H_0} h^2 = 0.0248$ without influence of other DM candidate.

Similar plots are shown in Figs. 3.5.11 and 12 with mass splitting $\Delta m = 25$ GeV. Figs. 3.5.9 and 11 depict similar nature and as we increase the value $\lambda_{\Phi X}$, the relation between m_X and m_{H_0} becomes more linear $m_X \sim m_{H_0}$. Although Fig. 3.5.11 is almost same as Fig. 3.5.9 when compared for different values of $\lambda_{\Phi X}$, due to the larger mass splitting contribution of $\Omega_{H_0} h^2$ to total dark matter relic density is smaller when compared with the case $\Delta m = 10$ GeV. This is shown in Fig. 3.5.12 and is obvious as for larger Δm values contribution in $\Omega_{H_0} h^2$ from co-annihilation channels are suppressed. However, still, there is sufficient contribution from the inert dark matter in total DM relic density. It can be observed that the value of $\Omega_{H_0} h^2 = 0.021$, for $m_{H_0} = 500$ GeV with $\Delta m = 25$ GeV when $\lambda_{\Phi X} = 0.3$. In the absence of the multicomponent scenario, the contribution of inert doublet candidate becomes $\Omega_{H_0} h^2 = 0.007$. Therefore, the extension of the inert doublet with a vector dark matter can

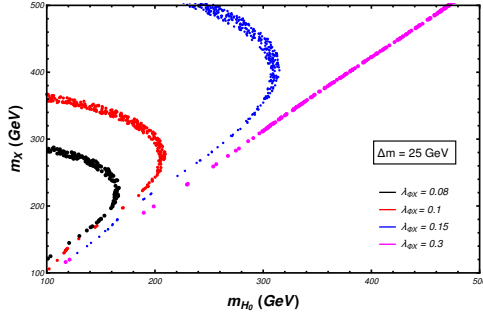


Figure 3.5.11: m_X vs m_{H_0} for $\Delta m = 25$ GeV as a function of $\lambda_{\Phi X}$.

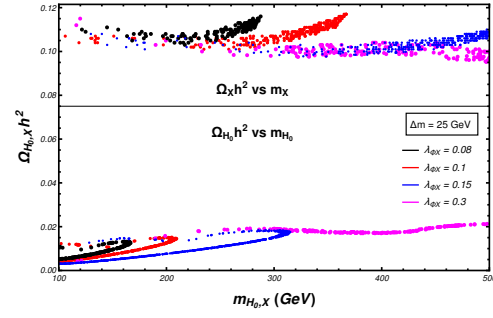


Figure 3.5.12: $\Omega_X h^2$ vs m_X , $\Omega_{H_0} h^2$ vs m_{H_0} for $\Delta m = 25$ GeV.

provide a successful multicomponent dark matter scenario where the contribution of the inert candidate can be enhanced considerably. However such properties are prominent for larger $\lambda_{\Phi X}$ values and when $m_X \sim m_{H_0}$. For smaller $\lambda_{\Phi X}$, the contribution of inert doublet increases up to a certain value satisfying the relation $m_X \sim m_{H_0}$ and then decrease with an increase in the mass of m_X . In such case, the inert doublet may have very small relic density (even when $m_X \geq m_{H_0}$) and most of the DM relic abundance is obtained from the other candidate X . Overall, from Figs. 3.5.9 and 11 we can conclude that $m_{H_0} \leq m_X$ region in m_{H_0} vs m_X plane is allowed in the present multicomponent model depending on the choice of $\lambda_{\Phi X}$.

Direct and Indirect detection of DM candidates: As mentioned earlier, from the study of vector dark matter [30], it is observed that although a vector dark matter can satisfy DM relic density with $\lambda_{HX} \sim 0.1$ for $m_X \sim 100 - 500$ GeV, it is ruled out by direct detection limits from XENON1T. For a smaller value of $\lambda_{HX} = 0.01$, it can be easily found that vector dark matter is in agreement with XENON1T limit for $m_X \geq 140$ GeV. However, for portal vector dark matter, relic density will be large as annihilation cross section becomes small. On the other hand in the two-component scenario, as shown in Fig. 3.5.4 when $m_X > m_{H_0}$, $\Omega_X h^2$ decreases due to new annihilation channels into an inert doublet. From Figs. 3.5.10–12, we also observe that contribution to the total DM relic density from vector dark matter in the multi-component framework is larger compared to the inert doublet. Therefore, conservatively assuming $\Omega_X h^2 \sim \Omega_{DM} h^2$, $\lambda_{HX} = 0.01$ allows us to study the low mass vector dark matter with $m_X > 140$ GeV which was excluded earlier.

Similar to the case of vector dark matter, the choice of $\lambda_L = 0.01$ for inert doublet dark matter is not arbitrary. Different collider searches for inert doublet dark matter also consider the limit on λ_L (also known as λ_{345}). Collider study of inert doublet including the direct detection limits from LUX excludes larger values of λ_L for $m_{H_0} \sim 100$ GeV and for $m_{H_0} = 500$ GeV, maximum allowed value of $\lambda_L = 0.4$ [166, 167]. Using the direct detection bound from XENON1T, the allowed limit on λ_L is further reduced. From [167] (see Fig. 7 in reference mentioned), it is found that for $m_{H_0} \sim 100$ GeV $\lambda_L \leq 0.01$ and for $m_{H_0} = 500$ GeV, $\lambda_L \leq 0.1$ is allowed. Therefore, in order to study inert doublet in the mass range $m_{H_0} \sim 100 - 500$ GeV, we consider a value of $\lambda_L = 0.01$.

As mentioned in the beginning of this Section 3.5, we considered $\lambda_L = \lambda_{HX} = 0.01$ and the direct detection cross section for both the candidates can be obtained from Eqs. (3.4.6) and (3.4.7). Therefore, the direct detection cross section is further reduced by the scaling r_i , $i = H_0, X$. We have found that, with the present choice of $\lambda_L = 0.01$, the regions plotted for H_0 in Figs. 3.5.1–3, 3.5.5, 3.5.7 and 3.5.9–12⁶ with $100 \text{ GeV} \leq m_{H_0} \leq 500 \text{ GeV}$ remains within the direct detection cross section limits obtained from XENON1T [29]. Similarly for $\lambda_{HX} = 0.01$ in case of the vector boson dark matter, we find that all the regions within the mass range $140 \text{ GeV} \leq m_X \leq 500 \text{ GeV}$ plotted in Figs. 3.5.6 and 3.5.8–12 is in agreement with the XENON1T bound. This limit is obtained by simply assuming $r_X \sim 1$ in Eq. (3.4.7) since in both Figs. 3.5.10 and 12 we observe that for masses $m_{H_0, X} \leq 200$ GeV, relic density of IDM

⁶There will not be any direct detection for Fig. 3.5.1 as $\lambda_L = \lambda_{HX} = 0$.

is small with respect to the vector dark matter. Therefore, all the allowed points with $m_X < 140$ GeV satisfying Eq. (3.3.3) in Figs. 3.5.9–12 are excluded by XENON1T bound on vector dark matter. This leaves an allowed parameter space in $m_{H_0} - m_X$ plane (shown in Figs. 3.5.9 and 11) satisfying the conditions i) $100 \text{ GeV} \leq m_{H_0} \leq 500 \text{ GeV}$, ii) $140 \text{ GeV} \leq m_X \leq 500 \text{ GeV}$ and iii) $m_{H_0} \leq m_X$ for fixed $\lambda_L, \lambda_{HX}, \Delta m$ considering $\lambda_{\Phi X}$ as the only variable. We use these limits for further studies involving collider searches for dark matter candidates in the next section. Of course one can avoid direct detection bound on vector dark matter by assuming a smaller value of λ_{HX} , thus consistent with XENON1T bound. But that would not affect the phenomenology of the model leaving $m_{H_0} \leq m_X$ as the only condition to restrain the model parameter space. It is to be noted that the parameter region discussed here, although safe from XENON-1T limits, can be probed by future DD experiments like XENONnT or DARWIN [12]. In fact, the low DM mass region (~ 100 GeV) can be excluded by DARWIN data in case of null result. This will further constrain λ_L and λ_{HX} to smaller values. But that will not affect the results because the study focuses on DM-DM conversion.

		Indirect detection		Direct detection	
m_X [GeV]	λ_{HX}	$\langle\sigma v\rangle_X$	$\langle\sigma v\rangle_{\text{Fermi-LAT}}$	σ_X	σ_{XENON1T}
		[cm ³ /s]		[cm ²]	
100	0.001	2.28×10^{-29}	3.0×10^{-26}	3.46×10^{-48}	1.5×10^{-46}
	0.01	2.27×10^{-27}		3.46×10^{-46}	
	0.02	9.0×10^{-27}		1.38×10^{-45}	
	0.03	2.0×10^{-26}		3.11×10^{-45}	
500	0.001	1.25×10^{-28}	1.0×10^{-25}	1.48×10^{-49}	4.0×10^{-46}
	0.01	1.37×10^{-26}		1.3×10^{-47}	
	0.1	9.1×10^{-30}		1.27×10^{-45}	
	1.0	5.3×10^{-26}		3.8×10^{-44}	

Table 3.5.1: Bound on λ_{HX} from direct and indirect detection. $m_{H_0} = 100$ GeV, $\lambda_{\Phi X}=0.1$, $\lambda_L=0.01$.

The bounds on the upper and lower limit of couplings λ_L and λ_{HX} , obtained from direct and indirect detection experiments, are given in Tables 3.5.1 and 3.5.2 respectively. It is self-explanatory that our choice of parameters are consistent with these constraints, with the comment that $m_X \leq 140$ GeV is excluded by direct detection limits, as already mentioned earlier which can also be verified from these Tables. While direct detection cross sections are purely Higgs mediated for both the DM candidates and therefore strongly depend on the Higgs portal couplings λ_L and λ_{HX} (see Eqs. (3.4.6) and (3.4.7) for expression), indirect detection cross sections may not have a similar dependence on the couplings, because, for H_0 , the dominant channel is annihilation into gauge pairs, which cannot be controlled by λ_L . However, indirect detection channels for X will depend on λ_{HX} since it is the only portal through which it can interact with SM. Note that $\sigma_{H_0} [\sigma_X]$ mentioned in the above tables are actually $\sigma_{H_0 \text{ SM} \rightarrow H_0 \text{ SM}} \times (\Omega_{H_0}/\Omega_{\text{DM}}) [\sigma_{X \text{ SM} \rightarrow X \text{ SM}} \times (\Omega_X/\Omega_{\text{DM}})]$, similarly, $\langle\sigma v\rangle_{H_0}$ and $\langle\sigma v\rangle_X$ are obtained by multiplying the total indirect annihilation cross sections for individual DM DM $\rightarrow$ SM SM multiplied by the fractions $(\Omega_{H_0}/\Omega_{\text{DM}})^2$ and $(\Omega_X/\Omega_{\text{DM}})^2$ respectively. The direct detection experimental bound is obtained from XENON1T [29] and indirect bound is taken from Fermi-LAT data [20]. It is to be noted that in the mass range $100 \text{ GeV} \leq m_{X,H_0} \leq 500 \text{ GeV}$ both the dark matter candidates annihilate into gauge bosons dominantly. The Fermi-LAT limit on the annihilation cross section is on DM annihilating directly into $b\bar{b}$. Both cross sections $\langle\sigma v\rangle_{H_0}$ and $\langle\sigma v\rangle_X$ are found to be much smaller to Fermi-LAT limit and therefore consistent with the bounds from indirect detection since secondary annihilation cross

m_{H_0} [GeV]	λ_L	Indirect detection		Direct detection	
		$\langle\sigma v\rangle_{H_0}$	$\langle\sigma v\rangle_{\text{Fermi-LAT}}$	σ_{H_0}	σ_{XENONIT}
		[cm ³ /s]		[cm ²]	
100	0.001	4.47×10^{-28}	3.0×10^{-26}	7.2×10^{-50}	1.5×10^{-46}
	0.01	4.14×10^{-28}		7.2×10^{-48}	
	0.1	2.18×10^{-28}		4.13×10^{-46}	
	1.0	2.24×10^{-29}		8.48×10^{-45}	
500	0.001	1.81×10^{-33}	1.0×10^{-25}	1.25×10^{-50}	4.0×10^{-46}
	0.01	1.7×10^{-29}		1.23×10^{-48}	
	0.1	9.1×10^{-30}		1.0×10^{-46}	
	1.0	4.26×10^{-31}		3.83×10^{-45}	

Table 3.5.2: Bound on λ_L from direct and indirect detection. $m_X = 600$ GeV, $\lambda_{\Phi X}=0.1$, $\lambda_{HX}=0.01$.

section into $b\bar{b}$ pair will be even smaller. Upper limits on $H_0 H_0 \rightarrow W^+ W^-$ cross section is calculated in Ref. [140] considering emission of gamma-ray in dwarf galaxies and also anti-proton flux. We found that the annihilation cross section of $H_0 H_0 \rightarrow W^+ W^-$ is smaller in the present multicomponent dark matter model compared to the observed limits in Ref. [140].

3.6 Dark Matter signatures at the LHC

In this section, we will discuss the issues about the collider signatures of the dark matter specific to our model. We specifically focus on the *dilepton plus missing transverse energy* channel ($2\ell + \cancel{E}_T$, $\ell = e, \mu$) in the present study. We choose at least 2 leading p_T leptons in the final states irrespective of their charge. Among the two DM candidates in our model, the vector DM, X , is produced only in the association with the SM Higgs boson, h , because it belongs to a separate $\mathbb{Z}'_2$ symmetry. As we are aware that the branching ratio of $h \rightarrow 2\ell$ is negligible, X does not give any significant signal cross section in the present LHC environment. Hence the collider signature of the DM here is entirely dictated by H_0 . Previously, Refs. [48, 124] studied the DM signature at the LHC for $2\ell + \cancel{E}_T$ channel for IDM alone whereas Refs. [59] studied the same signature in the context of multiparticle DM model consisting of IDM plus scalar. In a separate study, the authors of Ref. [119] studied the multilepton channel with more than 2 leptons in the final state. Here we will not address such channels as the signal cross section will decrease significantly for them.

For this study, we first generate our model file which can be used in a Monte Carlo event generator. We build our model file using FeynRules [23]. This model file is then used for generation of events through MadGraph5 [26]. We further use Pythia 8 [215] for showering, fragmentation and hadronization to get .HEPMC files. These .HEPMC files are then processed for fast simulation in detector studies through Delphes [134]. We generate events for the LHC at the CM energy $\sqrt{S} = 13$ TeV. We used the dynamic factorisation and renormalisation scale for the signal as well as the background events.

For the generation of parton-level events we apply minimum or maximum cuts on the transverse momentum p_T and rapidities η of light jets, b -jets, leptons, photons and missing transverse momentum $\cancel{E}_T$. Also, distance cuts between all possible final objects in the rapidity-azimuthal plane are applied, with the distance between two objects i and j defined as $\Delta R_{ij} = \sqrt{(\phi_i - \phi_j)^2 + (\eta_i - \eta_j)^2}$, where ϕ_i and η_i are the azimuthal angle and rapidity of the object i , respectively.

The preliminary selection cuts are:

- $p_T > 10$ and $|\eta| < 5$ for all *non-b*-jets, photons and leptons, and
- $\Delta R_{ij} > 0.4$ between all possible jets and leptons or photons.

After this, the .LHE files obtained through parton level events are showered with final-state radiation (FSR) with Pythia 8 where initial state radiation (ISR) and multiple interactions are switched off and fragmentation/hadronization is allowed.

We studied the following signal processes for the $2\ell + \cancel{E}_T$ final state.

Sig1: $pp \rightarrow A_0 H_0$, with A_0 further decaying through the channel, $A_0 \rightarrow \ell^+ \ell^- H_0$.

Sig2: $pp \rightarrow H^+ H^-$. $H^\pm$ can further decay into a pair leptons in following three ways:

- Both the charged Higgs bosons decay as $H^+(H^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)A_0$, followed by $A_0 \rightarrow \ell^+ \ell^- H_0$;
- Either one of the charged Higgs bosons decays as $H^+(H^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)A_0$, followed by $A_0 \rightarrow \ell^+ \ell^- H_0$, whereas the other one decaying through $H^+(H^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)H_0$.
- Both the charged Higgs bosons decay as $H^+(H^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)H_0$.

Sig3: $pp \rightarrow H^\pm H_0$. $H^\pm$ can further decay into a pair leptons as $H^+(H^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)A_0$, followed by $A_0 \rightarrow \ell^+ \ell^- H_0$

	m_{H_0}	m_{A_0}	m_{H^+}	m_X	λ_L	λ_{HX}	$\lambda_{\Phi X}$	$\Omega_{H_0} h^2$
	[GeV]							
BP1	150	160	160.1	205	0.01	0.01	0.07	0.01390
IDM	150	160	160.1	205	0.01			0.00468
BP2	100	125	125.1	155	0.01	0.01	0.05	0.00901
IDM	100	125	125.1	155	0.01			0.00197

Table 3.6.1: Viable benchmark points used for collider study.

We neglected any other signal process as they are negligible in comparison to the above-mentioned ones for our benchmark points given in Table 3.6.1. We have chosen these benchmark points based on the analysis of the relic density described in the previous sections. In Table 3.6.1 we have also given a comparison between IDM and our model for the same benchmark points.

The major background at the LHC for the $2\ell + \cancel{E}_T$ final state processes are as follows

Bkg1: $pp \rightarrow t\bar{t}$, with the top (anti-)quark decaying into the leptonic channel, $t(\bar{t}) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)b(\bar{b})$.

Bkg2: $pp \rightarrow W^+ W^-$. $W^\pm$ further decays via leptonic channel as $W^+(W^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)$.

Bkg3: $pp \rightarrow W^\pm Z(\gamma^*)$, followed by $W^+(W^-) \rightarrow \ell^+(\ell^-)\nu_\ell(\bar{\nu}_\ell)$, and Z/γ^* decays into leptonic channel, $Z(\gamma^*) \rightarrow \ell^+ \ell^-$.

Bkg4: $pp \rightarrow ZZ(\gamma^*)$, followed by leptonic decays $Z \rightarrow \nu_\ell \bar{\nu}_\ell$ and $Z(\gamma^*) \rightarrow \ell^+ \ell^-$.

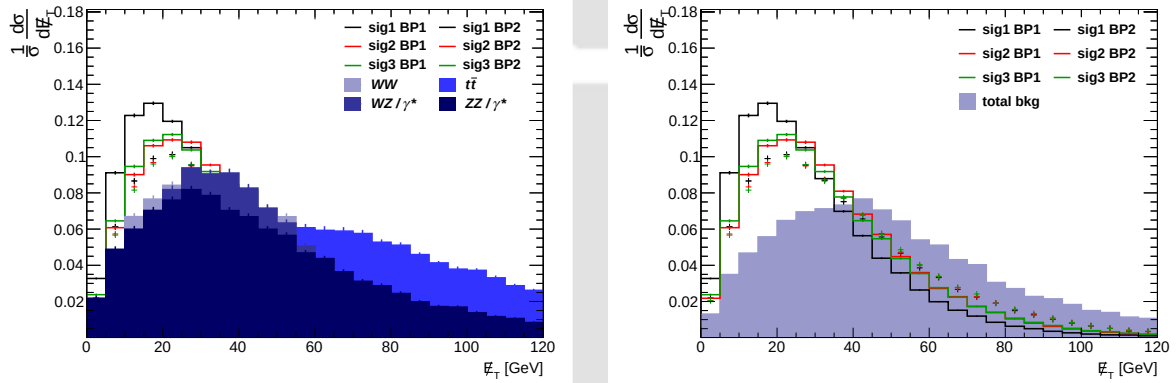
Table 3.6.2 shows the cross sections of the signal processes for the above-mentioned benchmark points along with their backgrounds.

As is obvious from Table 3.6.2 that the cross section is greater in the low mass region than the higher one, we can naïvely assume that the lower mass region holds much more promise in the search of new BSM signal. In Fig. 3.6.1 we showed the $\cancel{E}_T$ distribution for our benchmark points which further strengthens our conviction. Please note the difference between the two plots of Fig. 3.6.1. In

Processes	Cross section [fb]		Processes	Cross section [fb]
	BP1	BP2		
Sig1	3.40	9.83	Bkg1	21269.0
Sig2	1.99	4.59	Bkg2	3127.47
Sig3	1.26	2.64	Bkg3	402.68
			Bkg4	272.43

Table 3.6.2: Cross sections of the signals and backgrounds.

Fig. 3.6.1a) we plotted the $\cancel{E}_T$ distribution with each individual background. Here we see that the signal is almost entirely overshadowed by the background which conforms with previous such studies. However Fig. 3.6.1b) shows that if we see the signal with the total normalised background we can clearly distinguish the signals from the backgrounds for a low mass region. Another important point in our study is that we refrain from a strong $\cancel{E}_T$ cut. Usually, the LHC searches for the dark matter was conducted in the context of SUSY theories. As a result, a strong $\cancel{E}_T$ cut ~ 100 GeV is applied to suppress the background which washes out the signal entirely in the low $\cancel{E}_T$ region. Here such a strong $\cancel{E}_T$ cut is not necessary for IDM alone, specifically for the benchmark points of our choice.

**a)** Distribution of $\cancel{E}_T$ signal with individual backgrounds.**b)** Distribution of $\cancel{E}_T$ signal with total background.**Figure 3.6.1:** Distribution of $\cancel{E}_T$ for the benchmark points given in Table 3.6.1. The solid histograms belong to the BP2 and the dashed histograms (with only the peak value in each bin) show that for BP1.

Finally, we can say that a large $\cancel{E}_T$ signal at the LHC where the dark matter can be produced along with other visible Standard Model particles, be it photons, jets or leptons, as in our case, can provide a mode of discovery of invisible sector. Although the measurement of $\cancel{E}_T$ relies heavily on the precise measurement of all the other visible particles in the collision, it is a powerful tool for observing dark matter. With the proposed upgrade in luminosity coming up, if we succeed in detecting dark matter at the LHC, it will further complement the searches from cosmic ray experiments and help us solve the riddle of the Universe.

3.7 Summary

In this chapter, we perform a detailed analysis of a multicomponent dark matter model. We consider a two-component dark matter model with inert Higgs doublet associated with a vector boson dark matter both odd under two different discrete symmetries. We explore the intermediate-mass regime of inert Higgs doublet in the range $100 - 500$ GeV and observe that in presence of the other DM candidate it is possible to enhance the relic density of inert dark matter H_0 considerably when compared with the

usual single component inert doublet model. We observe that the allowed mass ranges of inert doublet and vector dark matter which satisfy total DM relic abundance depend significantly on the coupling between inert Higgs doublet and vector dark matter $\lambda_{\Phi X}$. We scan over a range of mass for both H_0 and X in the range $100 \text{ GeV} \leq m_{H_0, X} \leq 500 \text{ GeV}$ for a specific value of $\lambda_L = \lambda_{HX} = 0.01$ and mass splitting Δm . We show that there exists a large allowed parameter space of the model in m_{H_0} vs m_X plane for different $\lambda_{\Phi X}$ values satisfying the condition $m_{H_0} \leq m_X$. Direct detection limits from XENON1T further restricts the model parameter space by excluding regions below $m_X < 140 \text{ GeV}$. However, this can be avoided by considering a smaller λ_{HX} that satisfies direct detection limits for $m_X \sim 100 \text{ GeV}$ making the total range $100 \leq m_X \leq 500 \text{ GeV}$ accessible. But this does not affect the phenomenology of multicomponent dark matter model. In a similar vein, we have shown that an allowed parameter space can be obtained in the above mentioned intermediate-mass range which obeys the indirect detection bound from Fermi-LAT. Since IDM already satisfies relic bound in the region $m_{H_0} \geq 550 \text{ GeV}$ onwards due to co-annihilation channels which become more effective in this mass range, we have not shown the effect of our analysis in this mass range here. But we checked the region for both DM masses up to 2 TeV, the features remain the same - only larger couplings are required in the higher mass range.

Apart from the phenomenology of the dark sector, we also explore the collider signature of the inert doublet dark matter. We observe a clear signal in the lower mass region of the new particles of the inert doublet model. However one should be careful in putting a hard $\cancel{E}_T$ cut. Otherwise, the signal will completely be washed away. The collider signatures of vector dark matter are not very significant and have not been pursued in this work. In conclusion, we show that the present model with inert doublet and a vector boson dark matter can successfully provide a multi-component dark matter scenario.

Appendix

3.A Vector DM annihilation cross sections

In this appendix, we give the analytical expressions of the most important conversion cross sections for vector DM.

$$\sigma_{XX \rightarrow H^+ H^-} = \frac{1}{144\pi s m_X^4} \frac{\sqrt{s - 4m_{H^\pm}^2}}{\sqrt{s - 4m_X^2}} (12m_X^4 - 4m_X^2 s + s^2) \times \left\{ 4 \frac{(\lambda_L v^2 + m_{H^\pm}^2 - m_{H_0}^2)^2 \lambda_{HX}^2}{(s - m_h^2)^2} - 4 \frac{(\lambda_L v^2 + m_{H^\pm}^2 - m_{H_0}^2) \lambda_{HX} \lambda_{\Phi X}}{(s - m_h^2)} + \lambda_{\Phi X}^2 \right\} \quad (3.A.1)$$

$$\sigma_{XX \rightarrow H_0 H_0} = \frac{1}{288\pi s m_X^4} \frac{\sqrt{s - 4m_{H_0}^2}}{\sqrt{s - 4m_X^2}} (12m_X^4 - 4m_X^2 s + s^2) \times \left\{ \frac{\lambda_L^2 \lambda_{HX}^2 v^4}{(s - m_h^2)^2} - 2 \frac{\lambda_L \lambda_{HX} \lambda_{\Phi X} v^2}{(s - m_h^2)} + \lambda_{\Phi X}^2 \right\} \quad (3.A.2)$$

$$\sigma_{XX \rightarrow A_0 A_0} = \frac{1}{288\pi s m_X^4} \frac{\sqrt{s - 4m_{A_0}^2}}{\sqrt{s - 4m_X^2}} (12m_X^4 - 4m_X^2 s + s^2) \times \left\{ 4 \frac{(\lambda_L v^2 + m_{A_0}^2 - m_{H_0}^2)^2 \lambda_{HX}^2}{(s - m_h^2)^2} - 4 \frac{(\lambda_L v^2 + m_{A_0}^2 - m_{H_0}^2) \lambda_{HX} \lambda_{\Phi X}}{(s - m_h^2)} + \lambda_{\Phi X}^2 \right\} \quad (3.A.3)$$

Part II

Coannihilation of scalar dark matter



In this part, we focus on a second possibility of reviving the shrunken parameter space of WIMPs through *coannihilation*. As already discussed in the Introduction Chapter 1, in the coannihilation phenomena [32, 42, 57, 105, 121, 128, 144, 148, 169, 170, 174, 184], DM annihilates with another nearly degenerate dark sector particle, and the chemical equilibrium between the annihilating particles ensures the substantial depletion of DM number density. As the direct detection channels are relevant only for $\text{DM DM} \rightarrow q\bar{q}$ interactions, the coannihilation works efficiently to keep the bounds from such measurements at bay. In this part, we take up the case of quintessential minimal WIMP paradigm of a scalar singlet as the DM candidate. Scalar singlet DM by itself is a very constrained model because of tight direct detection bounds. We discussed two scenarios in the two chapters, the latter being the logical continuation of the first. Both of these help us in escaping such bound and opening up new possibilities of both rich phenomenologies in relic density analysis and search prospects in the controlled laboratory environment.



Signatures of leptophilic scalar dark matter in coannihilation scenario

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4.1 Motivation

In this chapter, we revisit the scenario of the scalar singlet dark matter with a pair of accompanying dark leptons which form a vectorlike Dirac fermionic doublet. Gauge invariance requires that the two components be degenerate at tree level and only a small mass splitting of the order of 300 MeV can be generated by the radiative corrections [81]. One can generate finite mass splitting at tree level in a gauge-invariant way by increasing the particle content in the model. In this article, we introduce an additional scalar triplet for this purpose whose vacuum expectation value (VEV) will be responsible for generating a small but sufficient mass splitting between the dark leptons. This mechanism is used to generate a finite mass of neutrinos in *type-II seesaw* models [155]. However, the triplet may or may not play any role in the phenomenology of the dark matter depending on the values of their masses. We assume them to be very heavy so that their effect is negligible apart from generating sufficient mass splitting between the dark leptons which will play an important role in our study. Now, this dark sector doublet couples with the SM leptonic doublet and the scalar singlet via a novel Yukawa interaction which is less explored in the literature. There are two distinct interesting features of this model: (i) Since the new dark sector fermions form a doublet, they will interact with the SM via gauge interaction as well as the new Yukawa coupling, which, in turn, will give rise to new annihilation channels, and (ii) such a Yukawa structure will open up new production channels for leptonic final states giving much enhancement in cross sections to search for dark matter in collider environments like the LHC through the said channel. Depending on the choice of parameters, here the DM annihilation can have three distinguishable stages, namely, pair annihilation, coannihilation, and mediator annihilation. Here, it is to be noted that coannihilation scenarios in WIMPs are mostly studied in the literature in the context of supersymmetric (SUSY) [32, 42, 127, 128] and coloured coannihilating particles [41, 67, 165]. Our model discusses a leptophilic context, and the coannihilation channels play an important role here due to the gauge interaction in the dark sector in addition to the new Yukawa coupling. This feature is significantly different from the cases

explored in the literature where the leptophilic Yukawa structure involves singlet dark sector partners with the DM candidate [79, 179].

As mentioned above, since the coannihilating partner¹ couples to the SM with a gauge as well as Yukawa coupling, the leptonic search channels get a boost in cross section from it, and it is only logical to probe the said channel for collider signatures. Moreover, the leptonic channel gives cleaner signals than the other channels. Still, the collider searches of dark matter are a very challenging prospect. Note that any leptophilic DM model like ours contributes to the calculation of muon $g - 2$. Very good agreement between the theoretical calculation and experimental measurements of muon $g - 2$, $\Delta a_\mu = a_\mu^{\text{Exp}} - a_\mu^{\text{SM}} = 268(63)(43) \times 10^{-11}$ [225], put a strong constraint on the new Yukawa couplings of the light SM leptons. However, there is no such bound for the production of τ leptons. Hence, it would be a good prospect to probe that channel for dark matter signatures in colliders. Our case is similar to the SUSY theories where stau is the coannihilating partner [105, 121, 128, 144, 169, 170, 174, 184]. In the SUSY scenario, the particle content is much larger than our minimalistic model, leading to more involved phenomenology. On the other hand, in a minimalistic model like ours, we have more handle to pinpoint the effects of coannihilation, and it is less probable to be lost amid other effects.

Despite the leptonic channel getting a boost, the cross section can still be smaller. So, to probe light leptonic channels effectively, one must follow sophisticated techniques to separate signals from the backgrounds. The *multivariate analysis* is one such prospect. We perform *boosted decision tree* (BDT) response to separate feeble light leptonic signals from the background in an effective manner. On the other hand, despite further enhancement in cross section, the τ leptons mostly decay into hadronic jets, resulting in difficulty in their reconstruction. We used τ -tagged jets from the detector simulation with 60% τ -tagging efficiency to perform the BDT response.

4.2 Model

As we described briefly in the previous section, we want a model where the dark sector will consist of one or more coannihilating partners in addition to the scalar singlet dark matter. So, we consider a vectorlike Dirac fermionic doublet $\Psi^T = (\psi^0, \psi^-)$ and a real scalar singlet ϕ in addition to the SM particles. To achieve the stability of the dark sector, both these new fields are odd under $\mathbb{Z}_2$ symmetry, whereas the SM fields are $\mathbb{Z}_2$ -even. Gauge invariance makes the masses of each field in the fermionic doublet degenerate. The difference between the masses can come only from radiative corrections which are of the order of 300 MeV. However, one can introduce extra fields in the model to increase the mass splitting. That is the reason why we introduce a $\mathbb{Z}_2$ -even scalar triplet in addition to the SM doublet scalar. However, this scalar triplet does not affect the phenomenology of dark matter in any way and serves the only purpose of tuning the mass splitting of new fermions. In Table 4.2.1, the quantum number assignments of the particles relevant to new interactions are shown.

	ℓ_L	e_R	H	Δ	Ψ	ϕ
$SU(2)_L$	2	1	2	3	2	1
$U(1)_Y$	$-1/2$	-1	$1/2$	1	$-1/2$	0
$\mathbb{Z}_2$	$+$	$+$	$+$	$+$	$-$	$-$

Table 4.2.1: Quantum number assignment of the relevant fields in our model. Electromagnetic charges are given by $Q = t^3 + Y$.

The real scalar singlet ϕ which is our DM candidate interacts with the SM via the Higgs-portal. As the other dark sector particles $(\psi^0, \psi^\pm)$ form an $SU(2)_L$ doublet, it interacts with the SM through gauge

¹Nonobservance of any new fermionic partner state in LEP2 which interacts with the SM leptons via Yukawa-type interaction puts a bound of $m_\psi > 104$ GeV on their masses. However, please note that here the new fermions interact strongly only with the τ lepton. Since LEP2 τ detection was not very precise, we must take this limit with a pinch of salt.

bosons. The coannihilating doublet Ψ couples with the SM leptonic doublet ℓ_L and the scalar singlet ϕ via a Yukawa interaction. This is novel in the sense that the widely used Yukawa structure in any new physics model consists of a scalar doublet which is the replica of the SM Yukawa interaction. Although this particular Yukawa structure is less explored in the literature, it fits the bill for all our requirements for this study. Although, as mentioned previously, the scalar triplet does not play any role in the DM phenomenology, we write the relevant terms in the Lagrangian nonetheless for the sake of completeness.

Hence, the resulting Lagrangian takes the form

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + \bar{\Psi} i \not{D} \Psi - \bar{\Psi} M_{\Psi} \Psi + \frac{1}{2} (\partial_{\mu} \phi)^2 - \frac{\mu_{\phi}^2}{2} \phi^2 - \frac{\lambda_{\phi}}{4} \phi^4 - \frac{\lambda_{h\phi}}{2} (H^{\dagger} H) \phi^2 \\ & + \text{tr}[(D^{\mu} \Delta)^{\dagger} (D_{\mu} \Delta)] - \mu_{\Delta}^2 \text{tr}[\Delta^{\dagger} \Delta] - [\mu H^T i \tau^2 \Delta^{\dagger} H + \text{h.c.}] \\ & - \lambda_{\Delta} [\text{tr}(\Delta^{\dagger} \Delta)]^2 + \text{tr}[(\Delta^{\dagger} \Delta)^2] - \lambda_{H\Delta} [H^{\dagger} H \text{tr}(\Delta^{\dagger} \Delta) + H^{\dagger} \Delta \Delta^{\dagger} H] \\ & - \lambda_{\phi\Delta} \phi^2 \text{tr}(\Delta^{\dagger} \Delta) - \frac{1}{\sqrt{2}} [y_{\Delta} \bar{\Psi}^c i \tau_2 \Delta \Psi + \text{h.c.}] - [y_{\alpha} (\bar{\ell}_{\alpha L} \Psi) \phi + \text{h.c.}], \end{aligned} \quad (4.2.1)$$

where $\mathcal{L}_{\text{SM}}$ is the SM Lagrangian, $M_{\Psi} = m_{\psi^0} = m_{\psi^-}$ is the bare mass term of the new fermionic doublet, and $D_{\mu} = \partial_{\mu} + i g_W t^a W^a + i g' Y B_{\mu}$ is the covariant derivative. The mass of the scalar singlet ϕ is given by $m_{\phi}^2 = \mu_{\phi}^2 + \lambda_{h\phi} v^2/2$, where the VEV $v = \sqrt{v_H^2 + v_{\Delta}^2} \approx 246$ GeV, v_H and v_{Δ} being the VEVs of the doublet and triplet scalar fields, respectively. The mass splitting between the dark lepton fields in the doublet comes out to be

$$\delta = |m_{\psi^0} - m_{\psi^+}| = y_{\Delta} v_{\Delta} \quad (4.2.2)$$

A discussion is in order here on the existing bounds that constrain the model parameters. The couplings which will play a significant role in the DM dynamics are the Higgs-portal coupling $\lambda_{h\phi}$ and the Yukawa couplings $y_{\ell}, \ell = e, \mu, \tau$. To put the bounds from the direct detection searches at bay, we have considered $\lambda_{h\phi} \lesssim 10^{-4}$, which also takes care of the invisible decay measurement. On the other hand, the muon $g - 2$ measurement puts a bound on the value of the Yukawa couplings of light leptons. Although it need not be so stringent, we still take a conservative choice of values at $y_e \sim y_{\mu} \lesssim 10^{-9}$. This leaves the third-generation Yukawa coupling y_{τ} to be the only one free from experimental constraints. However, one must note that to keep our model in the perturbative regime, we must have $y_{\tau} \leq 4\pi$. The measurement of the ρ parameter [225] puts a bound on the value of the triplet VEV: $v_{\Delta} \sim 3$ GeV. So for a safe choice, we have taken the value of the mass splitting $\delta \leq 10$ GeV.

4.3 Dark matter phenomenology

4.3.1 Formalism

In the proposed model, DM number changing processes are (i) pair annihilation ($\phi\phi \rightarrow \text{SM SM}$), (ii) coannihilation ($\phi\psi^{\pm 0} \rightarrow \text{SM SM}$) and (iii) mediator annihilation ($\psi^{\pm 0} \psi^{\mp 0} \rightarrow \text{SM SM}$). The choice of parameters will determine the relative contribution of these processes towards the relic density as we discuss in the following sections. In agreement with the common assumption of thermal freeze-out, the dark sector particles are in equilibrium with the thermal bath in the early Universe. At the same time, they are also in chemical equilibrium with each other, due to substantial interaction strength between themselves. Keeping all these in mind, one can write the Boltzmann equation as follows [148]

$$\frac{dn}{dt} = -3Hn - \langle \sigma_{\text{eff}} v \rangle (n^2 - n_{\text{eq}}^2), \quad (4.3.1)$$

where n and n_{eq} are the DM number density and the equilibrium number density respectively. Now, the effective velocity averaged annihilation cross section $\langle \sigma_{\text{eff}} v \rangle$ specific to this model can be written as

$$\begin{aligned} \langle \sigma_{\text{eff}} v \rangle = & \frac{1}{[g_{\phi} + \bar{g}_{\psi^0} + \bar{g}_{\psi^{\pm}}]^2} [g_{\phi}^2 \langle \sigma_{\phi\phi \rightarrow \text{SM SM}} v \rangle + g_{\phi} \bar{g}_{\psi^0} \langle \sigma_{\phi\psi^0 \rightarrow \text{SM SM}} v \rangle + g_{\phi} \bar{g}_{\psi^{\pm}} \langle \sigma_{\phi\psi^{\pm} \rightarrow \text{SM SM}} v \rangle \\ & + \bar{g}_{\psi^0}^2 \langle \sigma_{\psi^0\psi^0 \rightarrow \text{SM SM}} v \rangle + \bar{g}_{\psi^{\pm}}^2 \langle \sigma_{\psi^{\pm}\psi^{\mp} \rightarrow \text{SM SM}} v \rangle + \bar{g}_{\psi^{\pm}} \bar{g}_{\psi^0} \langle \sigma_{\psi^{\pm}\psi^0 \rightarrow \text{SM SM}} v \rangle] \end{aligned} \quad (4.3.2)$$

where

$$\begin{aligned}\bar{g}_{\psi^0} &= g_{\psi^0} (1 + \Delta m_0)^{3/2} \exp[-x \Delta m_0], \\ \bar{g}_{\psi^\pm} &= g_{\psi^\pm} (1 + \Delta m_{\text{ch}})^{3/2} \exp[-x \Delta m_{\text{ch}}].\end{aligned}\quad (4.3.3)$$

In the expressions above, $g_\phi = 1$, $g_{\psi^0} = g_{\psi^\pm} = 2$ are the internal degrees of freedom and $x = m_\phi/T$. Δm 's are dimensionless mass splitting parameters defined as

$$\begin{aligned}\Delta m_0 &= (m_{\psi^0} - m_\phi)/m_\phi, \\ \Delta m_{\text{ch}} &= (m_{\psi^\pm} - m_\phi)/m_\phi.\end{aligned}\quad (4.3.4)$$

As previously mentioned, the pair annihilation and coannihilation channels predominantly control the DM freeze-out. The mass splitting between ϕ and other dark sector particles and the Yukawa couplings mainly determines the contribution of these processes towards the total DM annihilation cross section. These mass splittings play a very important role, especially for the coannihilation and mediator-annihilation processes, as the Boltzmann factor in Eq. (4.3.2) gives rise to a significantly increased annihilation cross section for small values of Δm 's.

Before we go into the details of the freeze-out mechanisms, let us discuss the parameters used in the analysis. Since the mass splitting parameters defined in Eq. (4.3.4) between DM and other dark sector particles play an important role in freeze-out of ϕ , we will use them as independent parameters along with DM mass. As discussed in the previous section, the only important Yukawa coupling here will be y_τ , which couples ϕ to the third-generation $SU(2)_L$ lepton doublet and the new fermionic doublet Ψ . In summary, we take the following values of the parameters throughout our analysis:

Free parameters: $m_\phi, \Delta m_0, \Delta m_{\text{ch}}, y_\tau,$
Fixed parameters: $\lambda_{h\phi} = 10^{-4}, y_e \sim y_\mu \sim 10^{-9}.$

(4.3.5)

4.3.2 Analysis and observations

Relic density

In addition to the Higgs-portal annihilation channels of scalar singlet DM, the present model introduces a Yukawa interaction between the dark sector particles and SM. Unlike the Higgs-DM quartic coupling ($\lambda_{h\phi}$), the new Yukawa coupling (y_τ) is unconstrained except for the perturbative limits. This provides an excellent tool to explain the relic density for a wide parameter space even with negligible $\lambda_{h\phi}$, which, in turn, alleviates the direct search bounds. We have performed DM analysis using micrOMEGAs [50].

Because of minuscule $\lambda_{h\phi}$, the Higgs-portal annihilation channels have a negligible contribution towards DM relic density, and, hence, we will focus on the newly introduced channels only. All these annihilation channels can be broadly classified into three categories:

- pair annihilation ($\phi\phi \rightarrow \text{SM SM}$) (Fig. 4.3.1),
- Coannihilation ($\phi\psi^{\pm 0} \rightarrow \text{SM SM}$) (Fig. 4.3.2), and
- Mediator annihilation ($\psi^{\pm 0}\psi^{\mp 0} \rightarrow \text{SM SM}$) (Fig. 4.3.3).

All these categories can coexist or supersede each other, depending on the choice of parameters. The coannihilation and mediator annihilation processes become efficient only for small mass splittings between the DM and the dark sector particles. This is due to the exponential factor sitting in the expression for the respective $\langle\sigma_{\text{eff}} v\rangle$ [see Eq. (4.3.2)]. Moreover, in this model, the cross sections of these processes are larger than the DM pair annihilation cross section, which is the precise reason why these processes significantly reduce DM relic density [75]. In the following analysis, we will see that, over the entire parameter space, these DM number changing processes supersede the pair annihilation contribution to the relic density by 1 or 2 orders whenever $\delta m = m_{\psi^\pm}(m_{\psi^0}) - m_\phi$ are small.

Similar to δm 's, the two Δm 's also account for the strength of the coannihilation and mediator annihilation channels for two heavier dark sector particles ψ^0 and $\psi^\pm$. Apart from this, the Yukawa coupling y_τ also plays a significant role. In this context, it is worth noting that, for the pair annihilation channels in Fig. 4.3.1, the cross section depends on y_τ^4 , while, for coannihilation channels, it is only a y_τ^2 dependence, and the mediator annihilation channels, being mostly gauge mediated, have very little dependence on y_τ .

Since all the above three categories of DM annihilation can coexist in the parameter space, it would be interesting to identify the limiting cases where the transition from one category to another is perceivable. It is worth mentioning here that the interaction channels between ϕ and two other dark sector particles are exactly similar (see Figs. 4.3.1 and 2), so in the degenerate mass limit, $m_{\psi^0} \sim m_{\psi^\pm}$ their contribution will be the same. However, here the extra scalar triplet allows a finite mass difference between $\psi^\pm$ and ψ^0 . This mass splitting helps to identify the dominant dark lepton in the number-changing DM coannihilation and mediator annihilation channels.

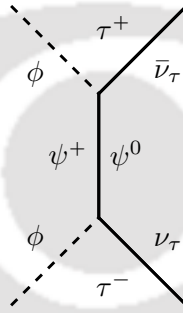


Figure 4.3.1: Feynman diagrams depicting the pair annihilation channels of ϕ

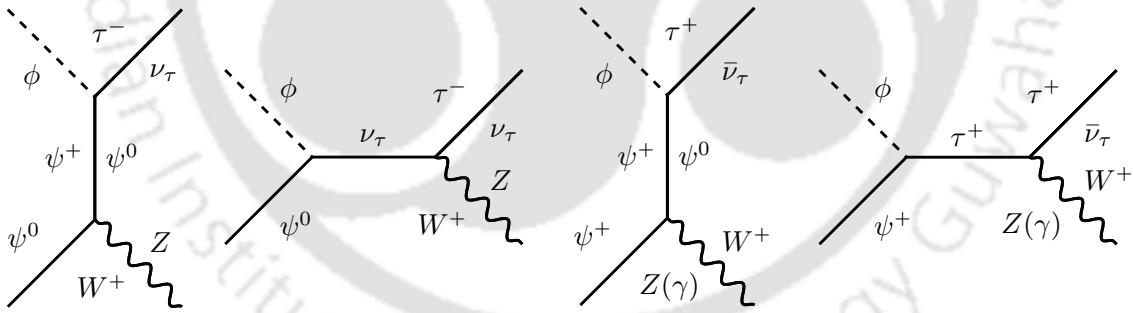


Figure 4.3.2: Feynman diagrams depicting the coannihilation channels of ϕ

In Fig. 4.3.4, the transition between the above-mentioned categories is depicted in the y_τ vs m_ϕ plane for some fixed values of $\delta m = m_{\psi^\pm}(m_{\psi^0}) - m_\phi$. For pair annihilation, $\langle \sigma_{\text{eff}} v \rangle$ has $m_{\psi^\pm}$ dependence only in the t -channel propagator, but, for coannihilation, $m_{\psi^\pm}$ appears in the propagator and the initial state along with the Boltzmann factor [Eq. (4.3.3)]. The functional dependence on the Boltzmann factor is even stronger for the mediator annihilation channels, which implies that for small mass splitting between the dark leptons and ϕ , mediator annihilation channels will deplete the DM number density most efficiently. For slightly larger splittings, the coannihilation channels take over. A large enough value of δm , however, makes the coannihilation processes negligible due to substantially large Boltzmann suppression, and so the pair annihilation predominantly dictates DM annihilation. This is clear from Fig. 4.3.4, where the blue line depicts the relic density allowed y_τ vs m_ϕ correlation for the mediator annihilation dominated channels. As required for this feature, $m_{\psi^\pm} - m_\phi$ is small ($\lesssim 10$ GeV). In magenta and red lines, however, coannihilation channels dominate due to comparatively larger

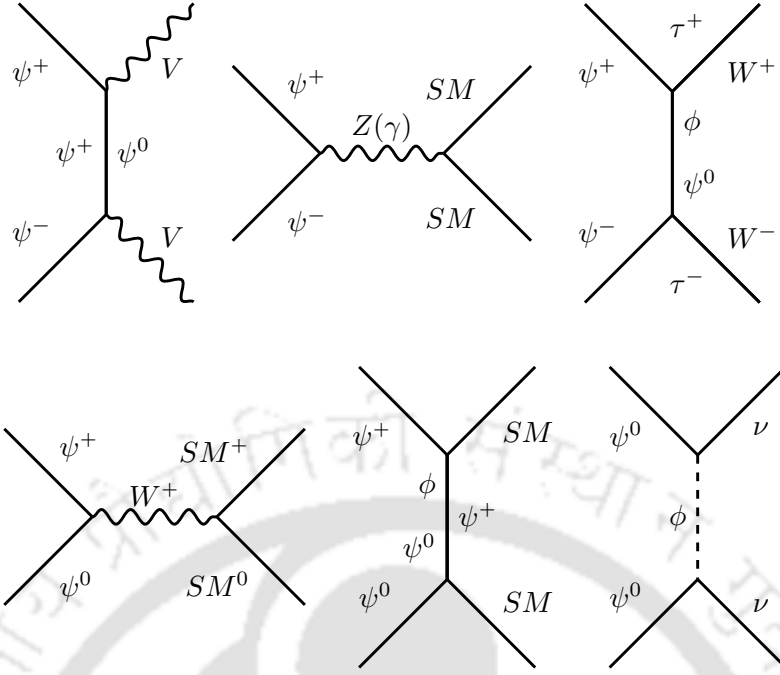


Figure 4.3.3: Feynman diagrams depicting the mediator annihilation channels

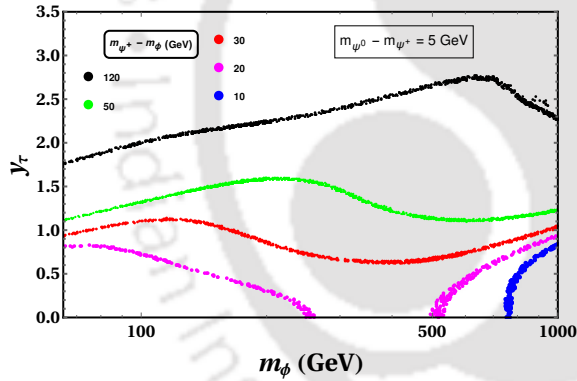


Figure 4.3.4: y_τ vs m_ϕ variation showing the transition from mediator annihilation to coannihilation and pair annihilation regime for different values of $m_{\psi^\pm} - m_\phi$ in GeV.

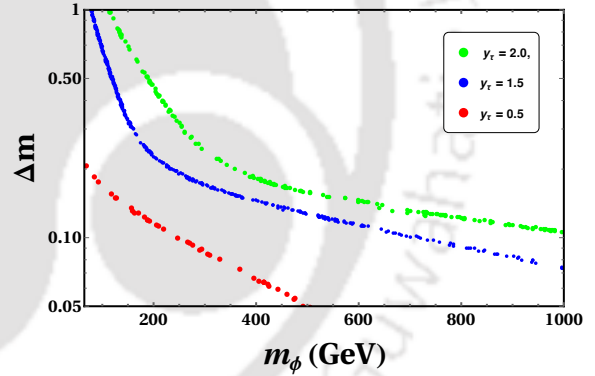


Figure 4.3.5: Δm ($= \Delta m_{\text{ch}} = \Delta m_0$) vs m_ϕ variation for fixed values of y_τ . All points satisfy relic density.

δm . It is to be noted here that $m_{\psi^0} - m_{\psi^\pm}$ is fixed at 5 GeV in this plot, which implies that $\phi - \psi^\pm$ coannihilation is stronger than the neutral dark lepton counterpart. As the DM-dark lepton splittings increase, for a fixed value of m_ϕ , $\langle \sigma_{\text{eff}} v \rangle$ becomes further Boltzmann suppressed, and, hence, larger y_τ is required to keep it within observed limits. For even larger δm , e.g., green and black lines, coannihilation contribution becomes more suppressed and pair annihilation becomes dominant. To obtain a sufficient annihilation cross section for the right relic, this implies a fairly large y_τ .

In Fig. 4.3.5, the variation of Δm ($\Delta m_{\text{ch}} = \Delta m_0 = \Delta m$) vs m_ϕ is plotted for all points satisfying the right relic. In the smaller m_ϕ region, the major share in relic density comes from $\phi\psi^{\pm(0)}$ coannihilation channels. Larger values of Δm cause Boltzmann suppression in $\langle \sigma_{\text{eff}} v \rangle$, which, in turn, is compensated by larger values of the coupling, as clearly seen in the plot. However, this feature is more prominent for smaller values of m_ϕ . This is because, for a fixed Δm , larger values of m_ϕ imply large $m_{\psi^{\pm(0)}} - m_\phi$, which rules out any substantial effect from the coannihilation channels. In fact, for larger m_ϕ , pair annihilation channels begin to dominate the total annihilation. Propagator suppression causes

y_τ to decrease towards higher m_ϕ , but here the decrease is at a much slower rate than the smaller m_ϕ regime.

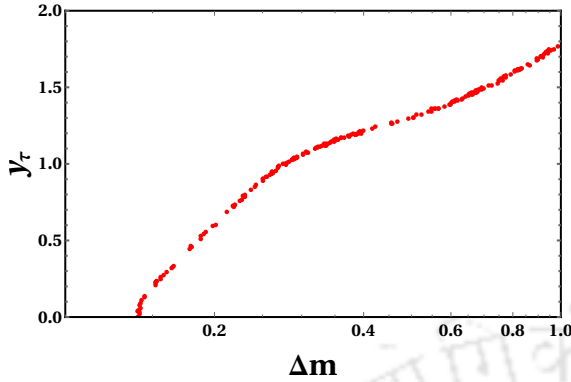


Figure 4.3.6: Variation of relic density allowed points in y_τ vs Δm ($= \Delta m_{\text{ch}} = \Delta m_0$) plane for $m_\phi=100$ GeV. The coupling gradually increases for larger Δm , compensating for larger Boltzmann suppression in $\langle \sigma_{\text{eff}} v \rangle$.

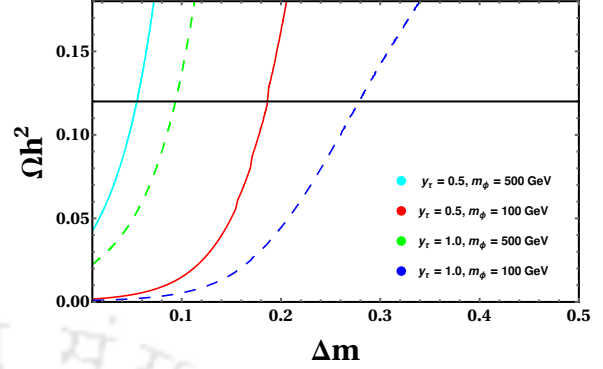


Figure 4.3.7: Variation of relic density with Δm ($= \Delta m_{\text{ch}} = \Delta m_0$) in the coannihilation regime for fixed values y_τ and m_ϕ . The black line represents the right relic density at $\Omega h^2 = 0.1215$.

As mentioned previously, for very small values of Δm 's, mediator annihilation or the freeze-out of ψ^0 and $\psi^\pm$ contributes to the relic density of ϕ . As observed from Eq. (4.3.2), the dependence on Δm 's is stronger in the Boltzmann factor of $\langle \sigma_{\text{eff}} v \rangle$ than coannihilation, and this leads to the fact that, for very small values of Δm 's, the mediator-driven annihilations almost entirely dominate the total DM annihilation. It is also worth noting that, being mostly gauge mediated, these channels substantially contribute to the DM annihilation even for very small values of y_τ . For our choice of parameters, we have observed that mediator annihilation is effective for $(m_{\psi^\pm(0)} - m_\phi) \lesssim 10$ GeV, and then the coannihilation processes take over. This feature is clear from Fig. 4.3.6, where we can see that there is no relic density allowed y_τ for $\Delta m \lesssim 0.1$ GeV. Beyond this range, as Δm increases, the required coupling also increases gradually to compensate for the Boltzmann suppression. m_ϕ is fixed at 100 GeV. Contributions from both the dark leptons are equal in the total annihilation cross section of DM since they are considered degenerate.

In Fig. 4.3.7, the variation of the relic density is plotted with Δm for some fixed Yukawa couplings and DM mass. As already argued, for a fixed m_ϕ , larger coupling corresponds to larger Δm due to Boltzmann suppression in $\langle \sigma_{\text{eff}} v \rangle$ as well as larger mass suppression of $\psi^{\pm(0)}$ in the t -channel propagator of the coannihilation channels. This explains the shift along the X axis from the red to the blue line where m_ϕ is 100 GeV and the Yukawa coupling y_τ varies from 0.5 to 1.0. We see the same trend for the cyan and green lines, but the amount of shift is relatively less, because, in this case, m_ϕ is larger (500 GeV), which automatically implies a fairly large splitting between $m_{\psi^\pm(0)}/m_{\psi^\pm}$ and m_ϕ , and, consequently, the coannihilation effect is not so prominent. We can argue that, for a fixed value of y_τ , larger DM mass obtains the correct relic density with a relatively smaller Δm ; hence, the red line with $m_\phi=100$ GeV shifts left towards the cyan line with the same y_τ but larger $m_\phi = 500$ GeV. The same logic applies to the shift between the blue and the green line. This trend also agrees with Fig. 4.3.5. As expected, very small values of Δm give an underabundance for the choice of parameters due to a fairly large increase in the Boltzmann factor of Eq. (4.3.3).

Now we observe the possibility of coexistence of all the possible annihilation regimes by varying both Δm_0 and Δm_{ch} along with the DM mass. Fig. 4.3.8 gives a correlation plot between Δm_0 and Δm_{ch} at different values of y_τ . For very small values of the coupling, mediator annihilation and coannihilation are dominant over pair-annihilation depending on Δm , because annihilation cross sections are proportional to y_τ^4 while for coannihilation it is only y_τ^2 and mediator annihilations are mostly gauge mediated. The green points, corresponding to $y_\tau = 0.0$, show that one can obtain the correct relic density only if one of the Δm 's is fairly small. Mediator annihilation is the only possibility here because some of the channels in Fig. 4.3.3 are y_τ independent. The red shading, on the other hand, corresponding to

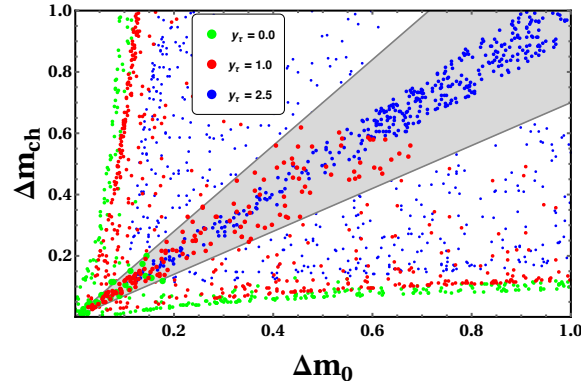


Figure 4.3.8: Correlation between Δm_0 and Δm_{ch} for different values of y_τ . m_ϕ varies from 65 GeV to 1 TeV. The three colours indicate three different couplings and distinguish the regions of parameter space where different DM number-changing processes dictate the relic density. The shaded region is allowed by the constraint on mass splitting between the dark leptons.

a larger coupling ($y_\tau=1.0$), shows some scattered points in the allowed region. This implies that, with the increase in the coupling, the coannihilation and pair annihilation channels become stronger, and, to maintain the right relic, the mediator annihilation channels are automatically suppressed. This is an artefact of the larger Δm values in the allowed region. The spread even increases for blue points, which corresponds to even larger $y_\tau = 2.5$. In this region, due to such large coupling, pair annihilation is the most dominant, and other number-changing processes are suppressed. Therefore, it becomes obvious that all three possible modes of annihilation can coexist in the present model for a wide parameter space where m_ϕ varies from a few GeV up to the TeV scale, and the coupling ranges from 0 to 3. However, the constraint on mass splitting between the dark leptons excludes the white region in the correlation plot. The introduction of a scalar triplet to generate the finite mass splitting plays the key role in relaxing the parameter space here because in the degenerate limit the allowed grey shaded region narrows down to the line along $\Delta m_0 = \Delta m_{\text{ch}}$. Nevertheless, even with the imposed constraint, the essential feature remains unchanged, where large Yukawa coupling and larger values of Δm favour pair annihilation while the smaller values facilitate coannihilation and mediator annihilation.

Direct and indirect detection

Direct search prospect — As known from the direct detection of scalar DM models, DM undergoes elastic scattering with detector nuclei through Higgs mediation. The spin-independent scattering cross section in our model is [47]

$$\sigma_{SI} = \frac{\lambda_{\phi h}^2}{16\pi m_h^4} f^2 \frac{m_N^4}{(m_\phi + m_N)^2} \quad (4.3.6)$$

where the form factor ($f \sim 0.3$) contains all the contributions from the nuclear matrix elements. Throughout the study, we have fixed the DM-Higgs coupling $\lambda_{\phi h}$ at 10^{-4} . This keeps σ_{SI} 2-4 orders below the present [29], as well as future [12] experimental bounds. The new physics Yukawa coupling y_τ being leptophilic plays no role in direct searches.

Indirect search prospect — The indirect detection experiments further constrain the DM velocity averaged cross section for relevant channels contributing to high-energy γ -ray flux in the Universe. In the context of our model, as far as these possibilities are concerned, due to DM-Higgs coupling $\lambda_{\phi h} = 10^{-4}$, $\langle\sigma v\rangle_{\gamma\gamma}$ and $\langle\sigma v\rangle_{b\bar{b}}$ contributions will be minuscule. However, the annihilation channels in Fig. 4.3.1 give rise to $\langle\sigma v\rangle_{\tau^+\tau^-}$ possibility.

In Fig. 4.3.9, $\langle\sigma v\rangle_{\tau\tau}$ is plotted against m_ϕ for two different values of Δm 's, whereas y_τ is varied in the colour bar. Both m_ϕ and y_τ are varied over the full range. To be specific, some values of y_τ are taken above our conservative choice for the perturbative limits to demonstrate the entire parameter space. The allowed limit ($y_\tau \lesssim 3.0$) is up to the green shade, whereas the purple region above is not

allowed by perturbativity. It is visible from Fig. 4.3.9b) that the large Δm 's considered here suggest larger propagator suppression for the relevant channels and, consequently, shift the parameter space downwards along the Y axis compared to Fig. 4.3.9a).

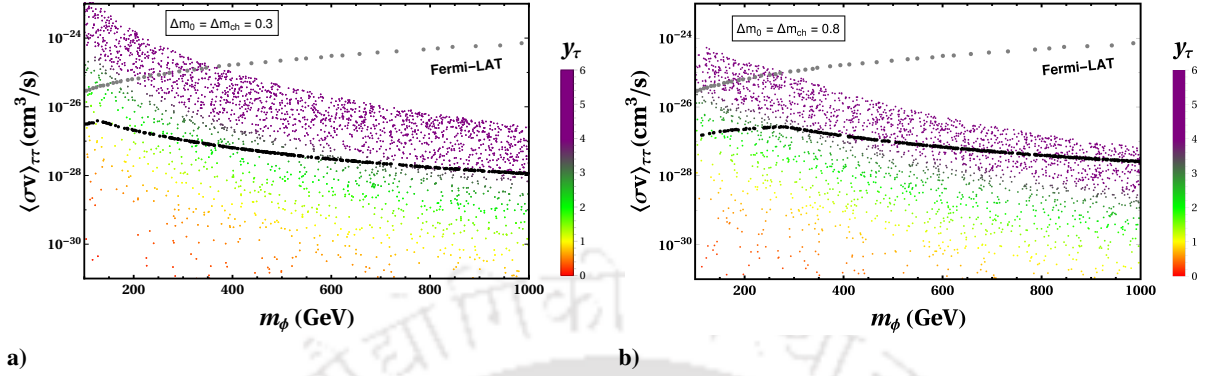


Figure 4.3.9: Plots indicating the parameter space and the constraints relevant to indirect detection. The top plot is for $\Delta m_0 = \Delta m_{ch} = 0.3$, whereas for the bottom plot $\Delta m_0 = \Delta m_{ch} = 0.8$. y_τ is varied in the colour bar, and our conservative choice for the perturbative limit is depicted by the region between red and green. The gray dots represent the latest bounds from $\tau^+\tau^-$ measurements observed in Fermi-LAT. The black dots represent the relic density allowed region.

The parameter space below the indirect search limits is depicted by the region below the grey dots, which is the latest experimental bound from Fermi-LAT data[15]. Because of the shift of the parameter region downwards along the Y axis for larger Δm 's as explained above, a larger region remains below the experimental limits for Fig. 4.3.9b), except a small portion towards smaller values of m_ϕ . However, it is seen that the relic density allowed region, depicted by the black dots, remains safely below the experimental limits in both the plots. For Fig. 4.3.9a), since the Δm 's are smaller, the relic density allowed region remains below the perturbative limits for y_τ for a large range of m_ϕ , whereas for Fig. 4.3.9b), a substantial portion remains above. This feature is also seen in Fig. 4.3.6. There, for a fixed value of DM mass ($m_\phi = 100$ GeV), the relic density allowed points correspond to $y_\tau \simeq 1.0$ for $\Delta m = 0.3$, whereas $y_\tau \simeq 1.7$ for $\Delta m = 0.8$.

4.4 Collider signatures

The challenges of discovering dark matter in colliders are manifold. They manifest themselves as *missing energy* (E_T^{miss}). Hence, the focus shifts entirely on the characteristics and precise measurements of associated production of visible particles. The charged multilepton channels are the most suitable to probe dark matter because of its clean signal, whereas QCD backgrounds overshadow the multijet channel, and it is very difficult to separate signals from the background. Since ours is a leptophilic model, these channels bear more significance than the others for our case. Having said so, please note that muon $g - 2$ constrains the new light leptonic couplings severely, whereas the coupling with the τ lepton remains unbounded, as we already mentioned previously. Here, we are going to study the collider signatures of DM through charged multilepton + E_T^{miss} channels. Our analysis will include both light charged leptons, since they give by far the cleanest signals, as well as τ leptons, as the unbounded couplings give a greater cross section than light leptons, paving the way to better analysis to see the effects of the new couplings.

Although τ lepton analysis poses more challenges, at the same time it unravels more unique features that can come in handy in the analysis for any collider like the LHC. The τ lepton is the charged lepton of the third generation and the heaviest among them. It is even heavier than most of the light quark mesons. As a result, τ leptons decay hadronically, which sets them apart from all other leptons. Because of the lepton number conserving weak interactions, the τ final states are always accompanied by one neutrino in the hadronic final states and two neutrinos in the leptonic final states. Since the neutrinos add

to the missing energy, the full τ energy cannot be measured. The leptonic decays of the τ are difficult to distinguish from prompt leptons in a $\ell + E_T^{miss}$ final state. Therefore, only the hadronically decaying τ 's are suitable for the collider signatures.

We have used FeynRules [23] to generate model files for our model. Events have been generated using MadGraph5 [26] and showered with Pythia 8 [215]. Finally, the detector simulation has been performed using Delphes [134]. We use the τ -tagged jets from Delphes and reconstruct them with the help of FastJet [83] using the anti- k_T algorithm. The separation ΔR of two adjacent τ jets is taken to be 0.4, and the τ -tagging efficiency is taken to be 60%. We carried out our analysis for the LHC at the CM energy $\sqrt{S} = 13$ TeV. We used the dynamic factorisation and renormalisation scale for the signal as well as the background events.

For the generation of parton level events, we apply minimum or maximum cuts on the transverse momenta p_T and rapidities η of light jets, b jets, leptons, photons, and missing transverse momentum. Also, distance cuts between all possible final objects in the rapidity-azimuthal plane are applied, with the distance between two objects i and j defined as $\Delta R_{ij} = \sqrt{(\phi_i - \phi_j)^2 + (\eta_i - \eta_j)^2}$, where ϕ_i and η_i are the azimuthal angle and rapidity of the object i , respectively.

The preliminary selection cuts used in the analysis are

- $p_T > 10$ GeV and $|\eta| < 2.5$ for all charged light leptons,
- $p_T > 20$ GeV and $|\eta| < 5$ for all *non-b* jets, and
- $\Delta R_{ij} > 0.4$ between all possible jets or leptons.

After this, the .LHE files obtained through parton level events are showered with *final state radiation* (FSR) with Pythia 8 where *initial state radiation* (ISR) and multiple interactions are switched off and fragmentation or hadronisation is allowed.

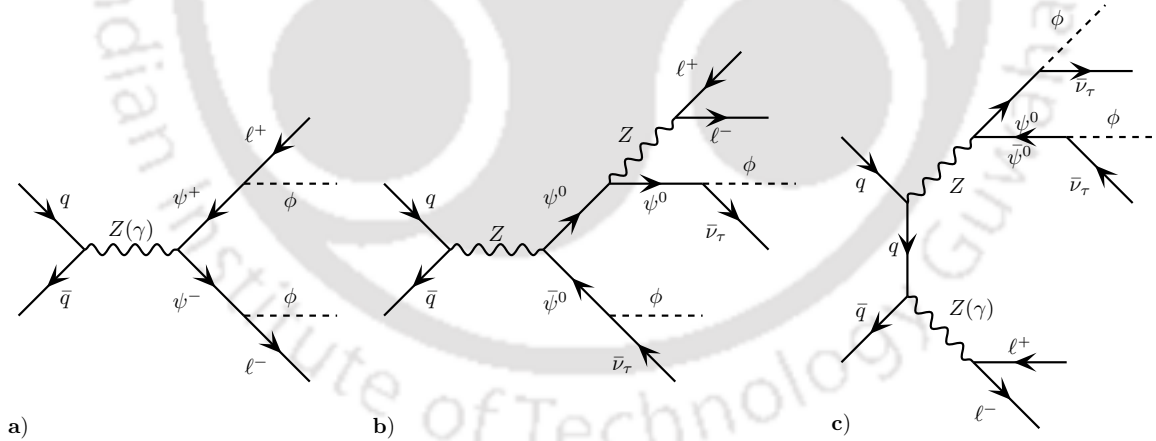


Figure 4.4.1: Feynman diagrams contributing to the dilepton channels.

The most important modes of production in dilepton channels are

- (1) $pp \rightarrow \ell^+ \ell^- 2\phi$;
- (2) $pp \rightarrow \ell^+ \ell^- \nu \bar{\nu} 2\phi$,

where ℓ stands for all three generations of charged leptons, namely, e, μ , and τ . Among the two classes of signal processes, the set (1) proceeds as follows:

- (i) $pp \rightarrow \psi^+ \psi^-$, followed by the decay of both ψ 's as $\psi^\pm \rightarrow \ell^\pm \phi$ [See Fig. 4.4.1a)] ,

The couplings which play a role in the above processes are only the ones involving light leptons and, hence, are very suppressed for the light dilepton channel as shown in Eq. (4.3.5). In the same vein, this channel will be dominant for the τ lepton analysis as a large value of y_τ is in effect.

On the other hand, the set (2) of processes mentioned above will proceed as

- (i) $pp \rightarrow \psi^0 \bar{\psi}^0$, followed by $\psi^0 \rightarrow \phi \nu$ and $\bar{\psi}^0 \rightarrow \ell^+ \ell^- \bar{\nu} \phi$ [See Fig. 4.4.1b)].
- (ii) $pp \rightarrow ZZ(\gamma^*)$, followed by $Z \rightarrow \nu \bar{\nu} 2\phi$ and $Z(\gamma^*) \rightarrow \ell^+ \ell^-$ [See Fig. 4.4.1c)].

As the couplings involved can also be either gauge couplings or that involving the τ lepton, both being considerably large, the set (2) of processes give us sufficient cross sections even for light dilepton channel to proceed with our analysis.

To highlight the features of our model clearly, we have selected the following benchmark points (see Table 4.4.1). The significance of the choice in benchmark points will be clear as we elaborate on our analysis in the following discussions.

	m_ϕ (GeV)	m_{ψ^0} (GeV)	$m_{\psi^\pm}$ (GeV)	y_τ
BP1	100	125	120	0.7
BP2	80	100	90	0.1

Table 4.4.1: Benchmark points used for the collider analysis.

For the benchmark points given in Table 4.4.1 we get the cross sections for the light dilepton channel as shown in Table 4.4.2.

Processes	Cross section (pb)	
	BP1	BP2
(2) $pp \rightarrow \ell^+ \ell^- \nu \bar{\nu} 2\phi$	232.17×10^{-6}	409.30×10^{-6}

Table 4.4.2: Cross sections of the light dilepton signal processes. The set (1) of subprocesses is highly suppressed for the coupling choices of Eq. (4.3.5), whereas the set (2) further involves gauge and τ lepton couplings and, hence, is dominant in this scenario.

Table 4.4.3 shows the signal cross sections for di- τ -jet channel. As mentioned previously, we can see the distinction between the cross sections of the process (1), which is substantially greater than that of the process (2) for this case. This is because, although the large value of y_τ and gauge couplings dictate both the processes, process (2) is suppressed by branchings and phase space.

Processes	Cross section (pb)	
	BP1	BP2
(1) $pp \rightarrow \tau^+ \tau^- 2\phi$	805.09×10^{-3}	786.33×10^{-3}
(2) $pp \rightarrow \tau^+ \tau^- \nu \bar{\nu} 2\phi$	2.72×10^{-3}	3.33×10^{-3}

Table 4.4.3: Cross sections of the signal processes for the di- τ -jet channel. Despite the large y_τ and gauge couplings, process (2) is suppressed due to more branchings and phase space.

The major backgrounds at the LHC for the light dilepton channel are as follows

Bkg1 : $pp \rightarrow t\bar{t}$, followed by the top (anti)quark decaying into the leptonic channel, $t(\bar{t}) \rightarrow \ell^\pm \nu(\bar{\nu}) b(\bar{b})$.

Bkg2 : $pp \rightarrow W^+W^-$. $W^\pm$ further decays via leptonic channel as $W^\pm \rightarrow \ell^\pm \nu(\bar{\nu})$.

Bkg3 : $pp \rightarrow W^\pm Z(\gamma^*)$, followed by $W^\pm \rightarrow \ell^\pm \nu(\bar{\nu})$, and Z/γ^* decays into leptonic channel, $Z(\gamma^*) \rightarrow \ell^+ \ell^-$.

Bkg4 : $pp \rightarrow ZZ(\gamma^*)$, followed by leptonic decays $Z \rightarrow \nu\bar{\nu}$ and $Z(\gamma^*) \rightarrow \ell^+ \ell^-$.

The major backgrounds at the LHC for the τ -jet channel will be similar as above with light lepton replaced by τ 's. Since τ 's can decay into hadronic channels we also have to consider light jets as backgrounds. In the following, we show all the backgrounds for this particular channel.

Bkg1 : $pp \rightarrow t\bar{t}$, followed by the top (anti)quark decaying into the τ -jet channel, $t(\bar{t}) \rightarrow \tau^\pm \nu(\bar{\nu}) b(\bar{b})$.

Bkg2 : $pp \rightarrow W^+W^-$. $W^\pm$ further decays via τ -jet channel as $W^\pm \rightarrow \tau^\pm \nu(\bar{\nu})$.

Bkg3 : $pp \rightarrow W^\pm Z(\gamma^*)$, followed by $W^\pm \rightarrow \tau^\pm \nu(\bar{\nu})$, and Z/γ^* decays into τ -jet channel, $Z(\gamma^*) \rightarrow \tau^+ \tau^- / 2j$.

Bkg4 : $pp \rightarrow ZZ(\gamma^*)$, followed by $Z \rightarrow \nu\bar{\nu}$ and jet decays $Z(\gamma^*) \rightarrow \tau^+ \tau^- / 2j$.

Table 4.4.4 shows the cross sections for the above backgrounds.

Processes	Cross section (pb)	
	Leptonic	Jet
Bkg1	21.27	5.31
Bkg2	3.13	781.43×10^{-3}
Bkg3	402.68×10^{-3}	14.75×10^3
Bkg4	272.43×10^{-3}	2.78×10^3

Table 4.4.4: Cross sections of the backgrounds.

Before getting involved in a more intricate analysis, we shall first discuss the kinematic distributions for this channel. The kinematic observables at our disposal are only the 4-momenta of the leptons and the missing energy E_T^{miss} . We order them according to the magnitude of their transverse momentum p_T . As a result, a *leading lepton* would always mean the leading- p_T lepton. We can also construct other observables from them, such as the invariant mass of the lepton pair. In a similar vein, we would construct the so-called transverse mass of the lepton- E_T^{miss} system. We shall call this quantity the missing transverse mass and define it for each lepton- E_T^{miss} system as

$$M_{T\ell}^{miss} = \sqrt{2 E_T^\ell E_T^{miss} (1 - \cos \Delta\phi_{\ell E_T^{miss}})}, \quad (4.4.1)$$

where $E_T^\ell = p_T^\ell$, which is the magnitude of the transverse momentum of a given lepton, and $\Delta\phi_{\ell E_T^{miss}}$ is the difference between the azimuthal angles of the lepton and missing transverse momentum. One can also construct similar observable for τ 's. The missing transverse mass plays an important role in distinguishing the massless invisible particles (such as neutrinos) from the massive ones (as is the case for our dark matter candidates), and, hence is very crucial for our analysis. With all these observables at our disposal, we show some of the distributions in Fig. 4.4.2 for the light dilepton channel and Fig. 4.4.3 for the di- τ -jet channel.

Before going further into the analysis, let us discuss the features of the benchmark points which we mentioned previously. It is clear from the distributions that the benchmark points 1 and 2 have very similar patterns and that the distributions are more populated in the lower region of each observable. It is also clear from these distributions that are very difficult to separate the signal from the background

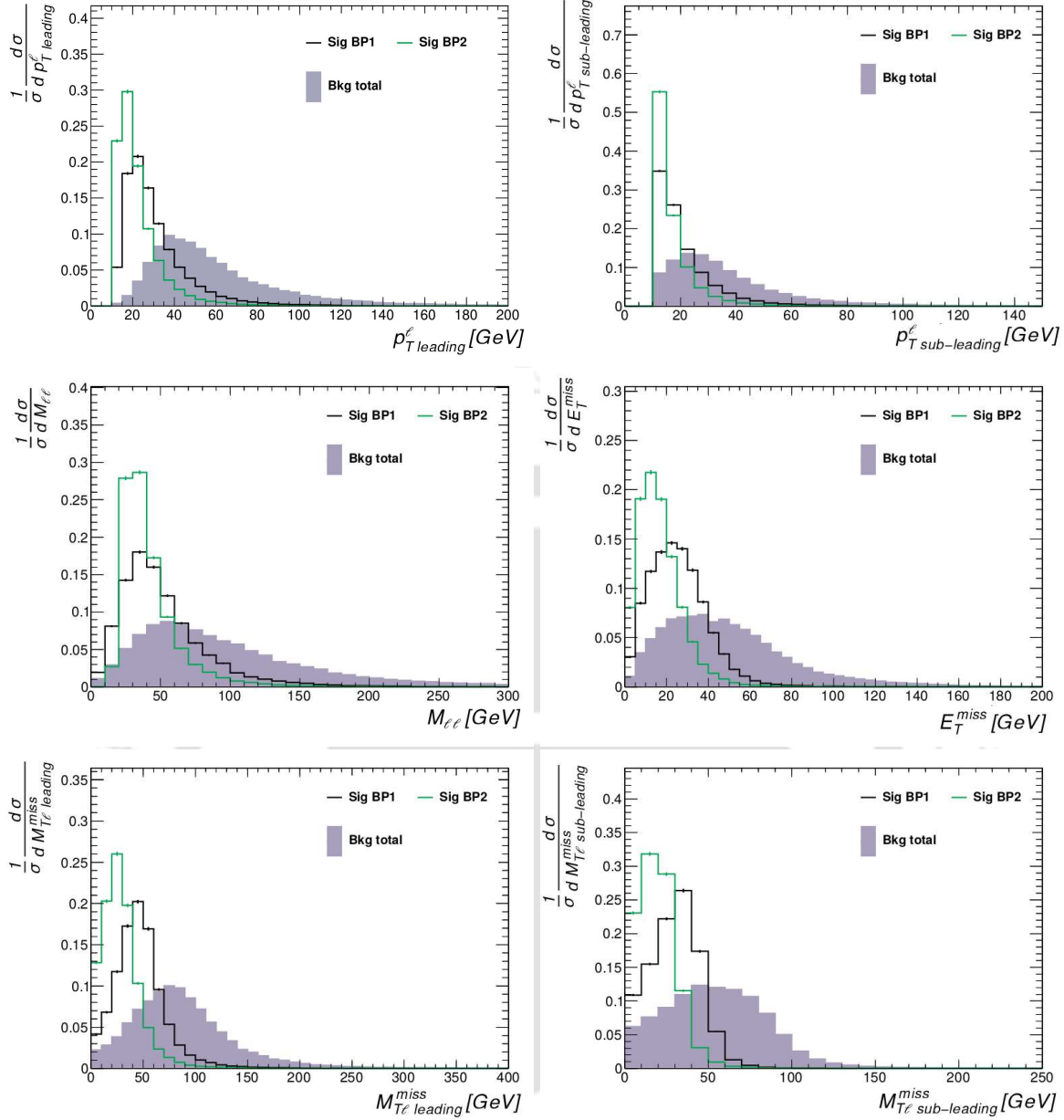


Figure 4.4.2: Distribution plots for the dilepton channel. The top panel shows p_T distributions of the leading and subleading p_T -ordered leptons and the lepton invariant mass. The bottom panel shows the distributions of the missing transverse energy and missing transverse masses with leading and subleading leptons.

by simple cut-flow analysis. We will lose too many signal events to the background leading to very low signal-to-background efficiency. That is the reason we apply the multivariate analysis which we will elaborate next.

In the next level of our study, we use *Toolkit for Multivariate Data Analysis* (TMVA) [164] in ROOT, to distinguish the signal events from the backgrounds efficiently. For this, we use the distributions of Fig. 4.4.2 (Fig. 4.4.3) and some other kinematic observables to train a BDT for the light dilepton (di- τ -jet) channel. The complete list of observables used to train BDT are as follows:

- p_T and η of the leading and subleading light leptons (τ -jets) and the invariant mass of the pair,
- missing transverse momentum E_T^{miss} ,
- missing transverse mass $M_{T\ell}^{\text{miss}}$ of the leading and subleading light leptons (τ jets),

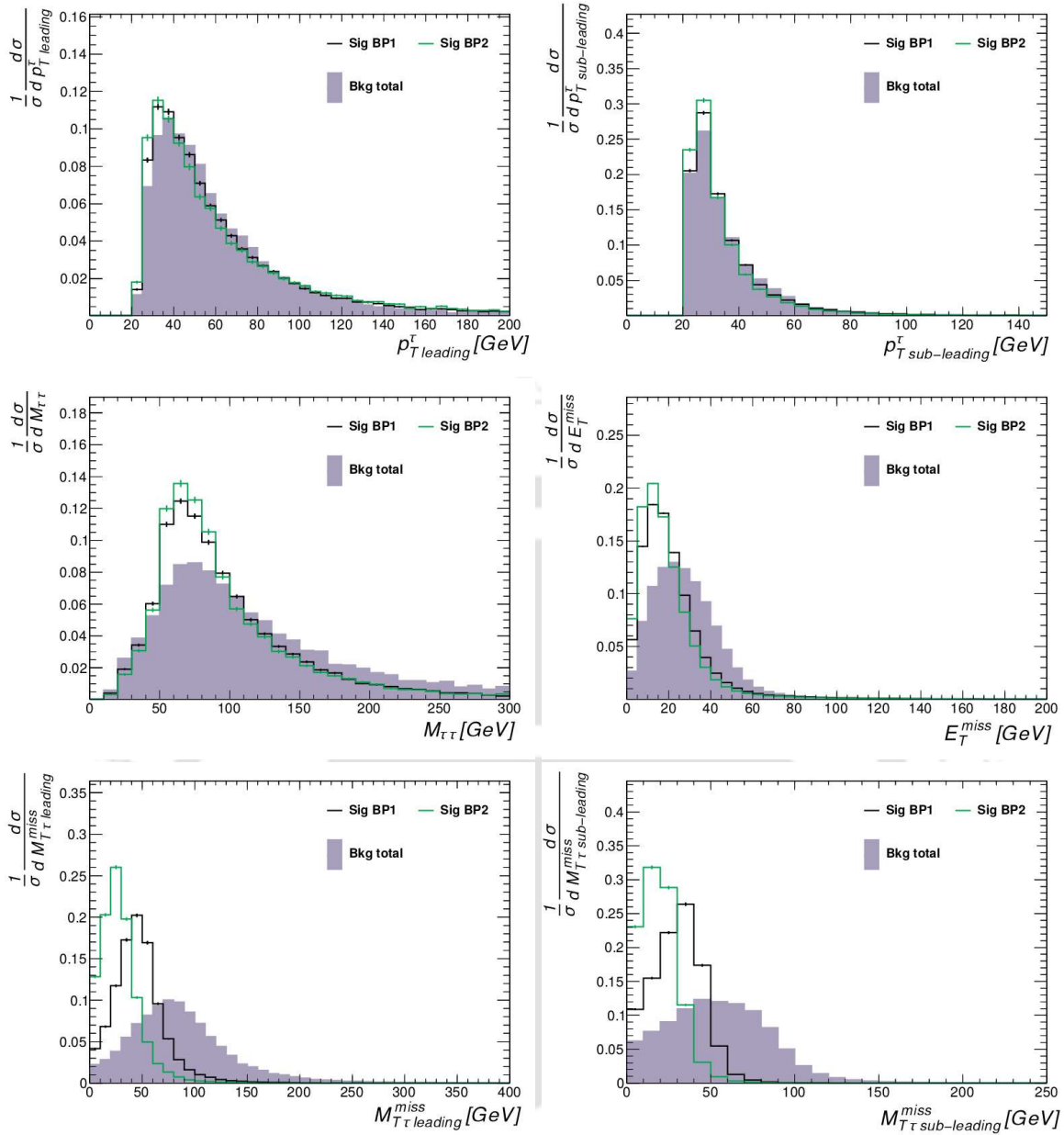


Figure 4.4.3: Distribution plots for the di- τ -jet channel. The top panel shows p_T distributions of the leading and subleading p_T -ordered τ jets and the τ jet invariant mass. The bottom panel shows the distributions of the missing transverse energy and missing transverse masses with leading and subleading τ jets.

- the difference of the azimuthal angles $\Delta\phi_{\ell E_T^{\text{miss}}}$ of the leading and subleading light leptons (τ -jets) with the missing transverse energy.

We use these distributions as discriminators to the BDT analysis. The discrimination of the signal and background can be improved further using proper cuts in addition to the preliminary selection cuts to the signal and/or background events. The resulting BDT response functions give us an estimate of the signal efficiency vs the rejection of the background.

Fig. 4.4.4 (Fig. 4.4.5) shows the BDT response curves (solid filled histograms are for the signal, and the hollow ones are for the background) for the light dilepton (di- τ -jet) channel for each benchmark points. Here we show two sets of BDT responses: (i) the solid purple and the hollow red ones are before any additional cuts with only taking into account the preliminary selection cuts, whereas (ii) the solid blue and the hollow black ones are after implementing a carefully chosen additional set of cuts to improve the distinguishability of the signal from the backgrounds. We will elaborate on the cuts chosen

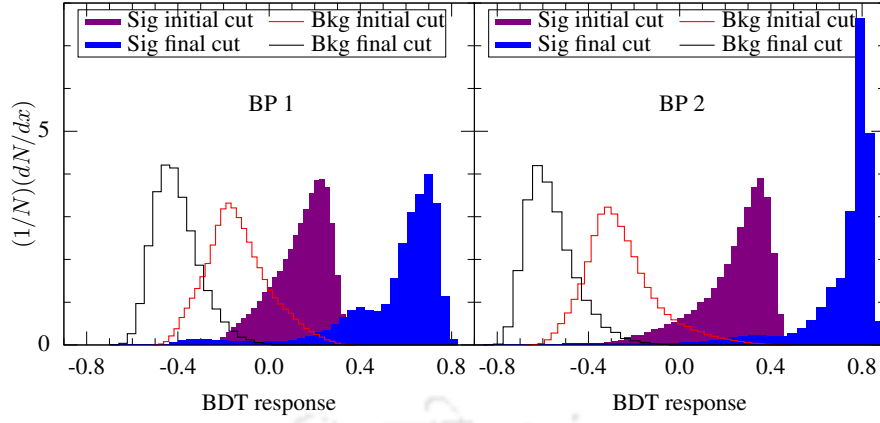


Figure 4.4.4: BDT response curves for the dilepton channel. The solid purple and the hollow red ones are before any additional cuts with only taking into account the preliminary selection cuts, whereas the solid blue and the hollow black ones are after implementing carefully chosen additional set of cuts.

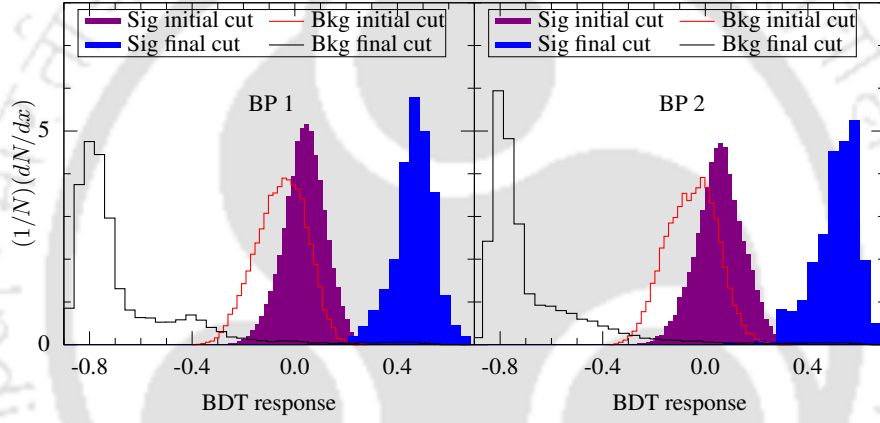


Figure 4.4.5: BDT response curves for the di- τ -jet channel. The solid purple and the hollow red ones are before any additional cuts with only taking into account the preliminary selection cuts, whereas the solid blue and the hollow black ones are after implementing a carefully chosen additional set of cuts.

later on.

We observe that the signal is separable from the background from the BDT response curves of Figs. 4.4.4 and 5 after the use of additional cuts. However, we have not yet quantified the improvement. For this, we draw the *receiver operating characteristic* (ROC) curves for each benchmark point using the gradual use of additional cuts. Fig. 4.4.6 (Fig. 4.4.7) shows the resulting curves for the signal efficiency vs the rejection of the background for the light dilepton (di- τ -jet) channel for each benchmark points. The area under each curve gives the quantitative estimate of the goodness of the separation of the signal from the backgrounds. In Table 4.4.5, we show the area under the ROC curve for each cut. The value of the cut is in addition to all the preceding cuts. We can see the improvement in the separation of the signal from the backgrounds from these numbers. The important point to be noted here is that, for the light dilepton case, the cuts were used only on the background events, leaving the signal events untouched, whereas, for the di- τ -jet case, the cuts were used only on the signal events. The reason for this can be understood as the population of events in the distribution plots of Figs. 4.4.2 and 3, respectively.

From the above discussions, plots, and numbers, we can see that, for our set of benchmark points, the DM signal for a leptophilic model with coannihilating dark lepton partners may not be separable from the SM backgrounds in a collider environment by the ordinary cut-flow analysis. However, it can be easily done with the help of a set of carefully chosen cuts in BDT analysis.

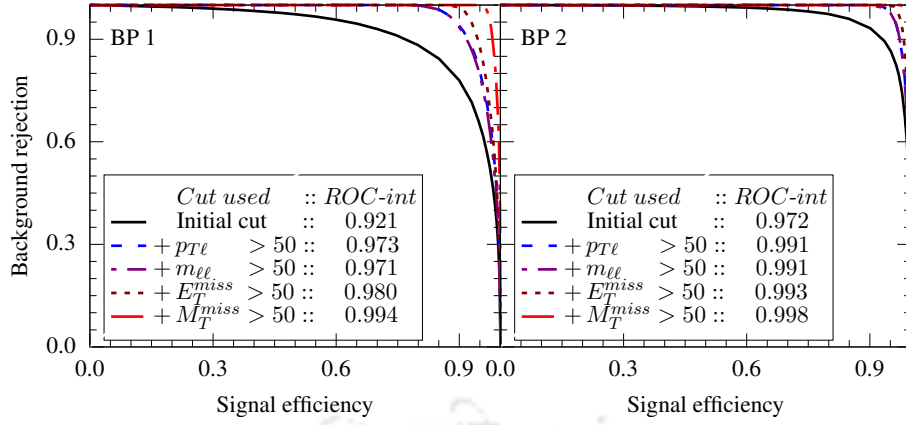


Figure 4.4.6: ROC curves for the dilepton channel. The area under the ROC curve for each cut (in GeV) is given in the inset. Each value is in addition to all the preceding cuts.

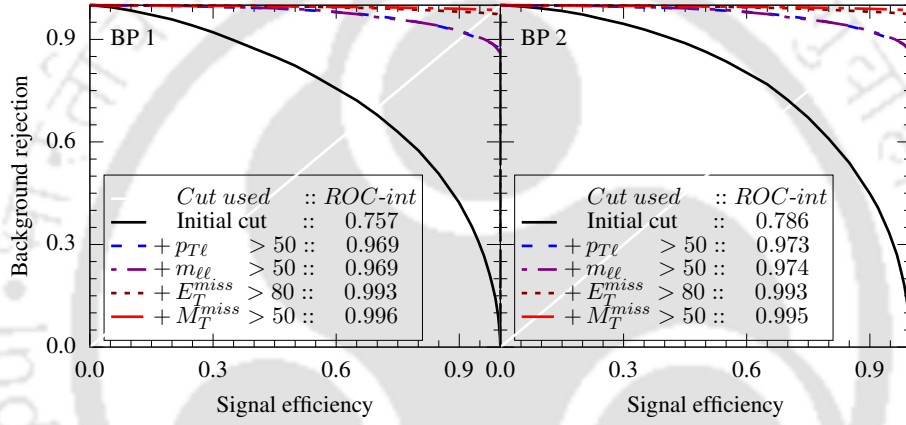


Figure 4.4.7: ROC curves for the di- τ -jet channel. The area under the ROC curve for each cut (in GeV) is given in the inset. Each value is in addition to all the preceding cuts.

dilepton channel			di- τ -jet channel		
Cuts	BP 1	BP 2	Cuts	BP 1	BP 2
Initial	0.921	0.972	Initial	0.757	0.786
$p_T^\ell > 50$ GeV	0.973	0.991	$p_T^\tau > 50$ GeV	0.969	0.973
$m_{\ell\ell} > 50$ GeV	0.971	0.991	$m_{\tau\tau} > 50$ GeV	0.969	0.974
$E_T^{\text{miss}} > 50$ GeV	0.980	0.993	$E_T^{\text{miss}} > 80$ GeV	0.993	0.993
$M_T^{\text{miss}} > 50$ GeV	0.994	0.998	$M_T^{\text{miss}} > 50$ GeV	0.996	0.995

Table 4.4.5: The area under the ROC curves in Figs. 4.4.6 and 7 for each cut. Each value is in addition to all the cuts above it. The efficiency of the cuts can be seen from the increasing values for each subsequent entry.

Apart from the prompt decay of dark partners, this model can also accommodate delayed decays in colliders leading long-lived particle (LLP) signatures which we discuss in the following. As mentioned previously, the gauge invariance mandates the degeneracy between the dark leptons. Only the loop effect generates a small mass splitting between them. Even after the introduction of a scalar triplet in our model, the mass splitting remains sufficient but small as can be seen from Eq. (4.2.2). The resulting

phase space will be automatically suppressed in the decays of such particles because of this small mass splitting. Such a scenario leads either of the dark leptons to be long-lived depending on their mass hierarchy. If the decay length $c\tau$ be greater than the detector radius, it is obvious that the neutral partner will give a missing energy signal, whereas the charged one will leave a stable ionisation track (HSCP) [2, 96, 158] in the detector.

Things become much more interesting once these particles decay inside the detector. For the charged particle, it will either give a disappearing (DT) [9, 137, 150] or a kinked track [38, 171].

Depending on the mass hierarchy, there can be several LLP possibilities. Fig. 4.4.8 shows such possibilities in our model. In Fig. 4.4.8a), ψ^+ decaying into a charged lepton and MET may give disappearing tracks or kinks in the tracker of the detector. The signature is similar to the charged Higgsino decaying into a neutral gravitino and SM leptons through off-shell decay of a W boson [171]. In our study, we are producing an on-shell ψ^0 in ψ^+ decay and that, in turn, decays into ϕ and ν_τ with 100% branching ratio. It is obvious that both the decay products of ψ^0 are charge-neutral, so the entire ψ^0 decay chain will give a missing energy signal in the colliders.

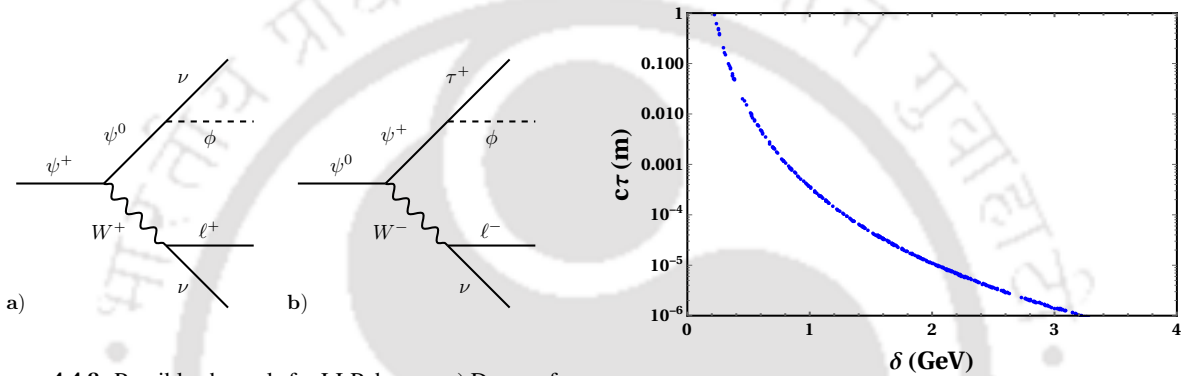


Figure 4.4.8: Possible channels for LLP decays. a) Decay of the charged dark lepton. b) Decay of the neutral dark lepton.

Figure 4.4.9: Variation of $c\tau$ as a function of the mass splitting between the dark leptons.

For nearly degenerate ψ^+ and ψ^0 , the emitted lepton will be soft, and the track of the mother particle can be identified as disappearing tracks. As we have already pointed out, radiative correction allows $\delta \sim 300$ MeV, and, from Fig. 4.4.9, it is clear that for this the in-flight decay of ψ^+ gives only DT of decay length around a few centimetres. The kink, however, may be observed in cases where the emitted lepton carries significant energy. One obvious way to address this issue is to increase the mass splitting (δ) between ψ^+ and ψ^0 . The introduction of the triplet in the particle spectrum becomes important in this context because δ can be varied up to 10 GeV. Fig. 4.4.9 shows the typical decay length ($c\tau$) vs δ variation for a typical value of m_{ψ^0} . The decay length varies around the sensitivity region of ATLAS for $\delta \lesssim 2$ GeV. For larger splittings, however, the decay becomes too prompt to be detected in the trackers. The kink angle, i.e., the difference of azimuthal angles of the charged parent and the daughter particle, can be measured as a function of the mass and 3-momentum of these particles and the polar angle of the parent particle with the beam axis. The expression, although model dependent, should be similar to Eq. (1) of Ref. [38]. However, the detailed kinematics specific to this model along with a robust collider simulation and proper background estimation is beyond the scope of this work, and we postpone it for future study.

If the mass hierarchy between ψ^0 and ψ^+ is reversed, then the neutral lepton decays late into two emerging visible particles and MET, the final states being a displaced jet coming from τ decay and a displaced lepton obtained via W decay [Fig. 4.4.8b)].

4.5 Summary

In this chapter, we have discussed a singlet scalar DM with a vectorlike fermionic doublet having the same dark symmetry. The minuscule Higgs-portal coupling with scalar DM keeps the direct detection

cross section below the experimental bound, which is an important handle in reviving the scenario of scalar singlet DM models. The new Yukawa coupling, on the other hand, which is irrelevant to direct search prospects, plays a vital role in dictating the relic density. We have shown that the model can provide a viable DM candidate through pair annihilation, coannihilation, and mediator annihilation channels over a wide parameter space ranging from GeV up to TeV scale. The transition from the pair annihilation to coannihilation regime is demonstrated, and the relevant limits of parameters are discussed. We have observed that coannihilation processes have a substantial contribution to relic density for a comparatively larger mass splitting between DM and the dark sector particles than what is usually discussed in the literature both in SUSY [37, 127, 130] and non-SUSY [93, 118, 179, 217] context. This may be attributed to the gauge couplings involved in these channels, which is a substantial contribution thanks to the dark fermion being a doublet. This is an artefact of the unconventional beyond the SM Yukawa structure considered in the proposed model. This arrangement, involving SM and dark sector lepton $SU(2)_L$ doublets and a scalar singlet, appropriately highlights the important features in the work.

Apart from the DM context, the gauge production of the fermionic doublet followed by decay to DM through the Yukawa coupling results in a substantially increased DM production at the colliders compared to scalar singlet scenarios. Using suitable kinematic observables in a BDT classifier, we separate the signal events from the backgrounds in an effective manner. We have shown that, with the use of proper cuts, we can achieve good results for both the light as well as τ leptonic channels.

In this work, the objective was to test the coannihilation scenario. This forced us to keep the mass splitting between the DM and the dark fermions small (~ 20 GeV). Also, as stated previously, the ρ parameter measurements further constrain the triplet vev, implying small splitting (~ 10 GeV) between the two dark fermions. This explains why larger mass was not possible to take. The DM mass could have been larger, but that would reduce the cross-section. Also, given the above constraints are applied, the profile of the distributions would not change. Similarly, larger coupling is also not favourable, since that would increase the pair annihilation contribution in the relic density and suppress coannihilation. However, one can in principle consider a larger mass split between the dark sector particles, as well as a larger couplings in the same model. In the next chapter, we have discussed this point in detail in a simple extension of this model, in the context of several multilepton + MET final states.

This model can also provide potential search prospects for long-lived particles because of the nearly degenerate or small mass splittings between the dark leptons. This can lead to suppressed phase space, and the delayed decay of these leptons can facilitate long-lived signatures (LLP) in the colliders, which is recently being given wide attention in the literature. In our study, we have indicated different possibilities of LLP signatures that may arise by tuning the relevant parameters.

One can interpret a limitation of the proposed model in the sense that, from the observed results in both dark matters and the collider analysis, there is no way to distinguish between the two dark leptons, although, one might assume on the contrary from Fig. 4.4.8 that it is possible from one-prong (for charged dark lepton) and two-prong (for neutral dark lepton) decays. But from Fig. 4.4.1 and the subsequent discussion, we see that we have to reconstruct a dark lepton from the visible final state SM particles; here, these are leptons. Since these two dark leptons have almost equal masses, it will be difficult to separate them, whereas, a larger splitting between the dark leptons can give more interesting signatures in the colliders. We are pursuing a possible solution to address these issues in the next chapter.

Implications of dark sector mixing on leptophilic scalar dark matter

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5.1 Motivation

In the previous chapter, we have shown how we can successfully evade the existing stringent constraints on scalar singlet DM by introducing a vector-like dark lepton doublet [91]. This simple addition to the scalar singlet DM extension of the SM opens up new vistas of possibilities, namely (i) new dark leptons interacting with the SM via gauge interactions introduces new annihilation channels, and (ii) novel Yukawa structure in the Lagrangian enhances cross sections facilitating its search in colliders.

However, the $SU(2)_L$ gauge invariance mandates the degeneracy of the two components of the doublet at the tree level. To lift this degeneracy, in Ref [91], we introduced a $\mathbb{Z}_2$ -even scalar triplet. The said scalar triplet otherwise does not play any role in the DM phenomenology, rendering the exercise ad hoc. Here we propose a better alternative to the earlier case, which is to add a $\mathbb{Z}_2$ -odd singlet fermion instead. In principle, the same charge dark fermions can mix among themselves to give physical eigenstates, which in turn, gives rise to the mass splitting between the pair. This mixing gives rise to rich phenomenological implication, unlike the previous case. Hence, this exercise is not limited to generating finite mass splitting between the components of the doublet. In other words, far from being an ad hoc addition to the model, this addition dictates the outcome of the processes through the mixing angle. Here we are interested in the implications of mixing between the charged dark sector fermions as well as the distinguishability of the pure and the mixed leptonic states from the observations.

To put the matter into perspective, we would like to mention that the mixing between a singlet and a doublet in the context of DM studies is nothing new. Previously, several authors used this structure to explain various phenomena. In supersymmetric (SUSY) theories, the bino can mix with the Higgsinos to give rise to a neutralino [35, 40, 97, 104]. Out from the realm of SUSY, this matter has been taken up by other authors to build up minimal DM models to address shrinking parameter space from relic density and direct detection measurements [46, 62, 85, 98, 108, 220]. In addition to the context of DM, this extension has successfully explained small values of neutrino masses and the related phenomenology

[60, 131, 206, 219]. See Ref. [58], for a recent review and an exhaustive list of relevant papers. Now, where we differ from all these works is that they took a mixed dark lepton state as a DM candidate, which limits the number of decay channels. Here, the DM candidate is a separate scalar singlet, which phenomenologically is the same as the quintessential minimal extension of the SM. The dark leptons add a new portal for pair annihilation of DM. Besides, there are coannihilation and mediator annihilation channels, and thus it enriches the DM dynamics. The mixing between the dark partners relaxes the existing experimental limits on the scalar singlet DM scenario. It also provides a handle to identify the viable parameter space dominated by dark sector partners of different isospins.

The probe of the DM candidates in the controlled collider environment is always a challenging task. Here we analyse the situation in the case of Large Hadron Collider (LHC). We focus on how the signatures of the mixing parameter from the kinematic distributions of the relevant observables can help us decipher the DM signals from our model. We will see that the addition of an extra singlet does not complicate the search strategies, but opens up new avenues. The mixing parameter gives an extra handle in tuning the kinematic distributions that will be easily accessible at the LHC. Apart from collider searches, we will see how does it affect the indirect search prospects of the DM.

5.2 Model description

We discussed briefly in the introduction the motivation for the choice of our model content. Here we shall develop from that motivation a detailed description of the model parameters. As mentioned previously, ours is a leptophilic model that can evade all the existing bounds on scalar singlet DM. For that, we consider dark-sector partner(s) which can either be a doublet [91] or a singlet fermion [179]. Previously [91] we studied the case of a dark lepton doublet partner which coannihilates with the singlet scalar DM candidate as well as acts as a portal, depending on the parameter space, which results in relaxing the constraints considerably. As the gauge invariance does not allow non-degenerate mass states for the doublet partners at the tree level, we had to introduce a new scale of electroweak symmetry breaking. We added a, otherwise redundant, scalar triplet to generate finite mass splitting between them. This mass splitting plays a significant role in determining the DM signatures in the collider. However, the measurement of the ρ parameter [225] constrains the value of the mass splitting ≤ 10 GeV. The motivation to make it arbitrary leads us to add a $\mathbb{Z}_2$ -odd singlet fermion instead of a $\mathbb{Z}_2$ -even scalar triplet in the particle spectrum. We shall see in this section how this singlet fermion plays a vital role in lifting the degeneracy of the doublet states depending on the parameter space.

Moreover, in Ref [91], we had pointed out that none of the observations was useful to distinguish between the two coannihilating particles. Here, in principle, the physical states of the dark leptons can be a mixed state between the same charge component of the doublet and the singlet. Thus it gives us the most general minimal scenario where we can meet all our demands. We want to distinguish between these two dark leptons from the observations of the DM in experimental searches. In this light, we will show later on, that the pair annihilations are more useful in the collider analysis and indirect searches. Whereas in the relic density scenario, the coannihilation channels can better exhibit the mixing effects. With all these motivations clear in our mind, we give below the parameter content of our model.

	ℓ_L	e_R	H	Ψ	ξ	ϕ
$SU(2)_L$	2	1	2	2	1	1
$U(1)_Y$	$-1/2$	-1	$1/2$	$-1/2$	-1	0
$\mathbb{Z}_2$	$+$	$+$	$+$	$-$	$-$	$-$

Table 5.2.1: Quantum number assignment of the relevant fields in our model. Electromagnetic charges are given by $Q = t^3 + Y$.

Table 5.2.1 shows the particle content and the quantum number assignments of our model. The $\mathbb{Z}_2$ -

odd dark sector contains a vector-like Dirac fermionic doublet, $\Psi^T = (\psi^0, \psi_1)$, a fermionic singlet ξ and a real scalar singlet ϕ . ϕ is our DM candidate which interacts with the SM Higgs via quartic as well as portal coupling and the $\mathbb{Z}_2$ symmetry renders stability to it. The fields Ψ and ξ couples with the SM doublets ℓ_L, H and the scalar singlet ϕ via three different kinds of gauge-invariant Yukawa interactions.

Hence the resulting Lagrangian takes the form

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + \bar{\Psi} (i \not{D} - M_{\Psi}) \Psi + \bar{\xi} (i \not{\partial} - m_{\xi}) \xi \\ & + \frac{1}{2} (\partial_{\mu} \phi)^2 - \frac{\mu_{\phi}^2}{2} \phi^2 - \frac{\lambda_{\phi}}{4} \phi^4 - \frac{\lambda_{h\phi}}{2} (H^{\dagger} H) \phi^2 \\ & - [y_j^D (\bar{\ell}_{jL} \Psi) \phi + y_j^S (\bar{e}_{jR} \xi) \phi + y (\bar{\Psi} H) \xi + \text{h.c.}] , \end{aligned} \quad (5.2.1)$$

where $\mathcal{L}_{\text{SM}}$ is the SM Lagrangian and a sum over the generation index j is implied. $M_{\Psi} = m_{\psi^0} = m_{\psi_1}$, is the degenerate bare mass term of the doublet and $D_{\mu} = \partial_{\mu} + i g_W t^a W^a + i g' Y B_{\mu}$ is the covariant derivative. The mass of the scalar singlet ϕ is given by

$$m_{\phi}^2 = \mu_{\phi}^2 + \lambda_{h\phi} \frac{v^2}{2}. \quad (5.2.2)$$

After electroweak symmetry breaking in Lagrangian (5.2.1), the Yukawa interaction involving the SM Higgs boson induces non-diagonal mass terms between the singlet and the same charge component of the doublet. As a result, the bare mass terms of the said fermionic degrees of freedom are augmented with the mixed terms as shown below.

$$\mathcal{L}_{\text{mass}} = m_{\psi_1} \bar{\psi}_1 \psi_1 + m_{\xi} \bar{\xi} \xi + \frac{y v}{\sqrt{2}} \bar{\psi}_1 \xi + \frac{y v}{\sqrt{2}} \bar{\xi} \psi_1 \quad (5.2.3)$$

The resulting mass matrix in the bare basis (ψ_1, ξ) is given by

$$M_{(\psi_1, \xi)} = \begin{pmatrix} m_{\psi_1} & \frac{y v}{\sqrt{2}} \\ \frac{y v}{\sqrt{2}} & m_{\xi} \end{pmatrix} \quad (5.2.4)$$

To diagonalise the above mass matrix, we introduce the physical basis (ψ, χ) through the orthogonal transformation

$$\begin{aligned} \psi &= c_{\alpha} \psi_1 + s_{\alpha} \xi , \\ \chi &= -s_{\alpha} \psi_1 + c_{\alpha} \xi ; \end{aligned} \quad (5.2.5)$$

where α is the mixing angle. Thus, the diagonal mass matrix of the physical mass terms turns out to be

$$M_{(\psi, \chi)} = \begin{pmatrix} m_{\psi} & 0 \\ 0 & m_{\chi} \end{pmatrix} = \begin{pmatrix} c_{\alpha} & s_{\alpha} \\ -s_{\alpha} & c_{\alpha} \end{pmatrix} \begin{pmatrix} m_{\psi_1} & \frac{y v}{\sqrt{2}} \\ \frac{y v}{\sqrt{2}} & m_{\xi} \end{pmatrix} \begin{pmatrix} c_{\alpha} & -s_{\alpha} \\ s_{\alpha} & c_{\alpha} \end{pmatrix}. \quad (5.2.6)$$

From the above, we get the physical masses m_{ψ}, m_{χ} and the mixing angle α in terms of the bare parameters

$$t_{2\alpha} = \frac{\sqrt{2} y v}{m_{\psi_1} - m_{\xi}}, \quad (5.2.7)$$

$$\begin{aligned} m_{\psi} &= \frac{1}{2} (m_{\psi_1} + m_{\xi}) + \frac{1}{2} \sqrt{(m_{\psi_1} - m_{\xi})^2 + 2 y^2 v^2}, \\ m_{\chi} &= \frac{1}{2} (m_{\psi_1} + m_{\xi}) - \frac{1}{2} \sqrt{(m_{\psi_1} - m_{\xi})^2 + 2 y^2 v^2}. \end{aligned} \quad (5.2.8)$$

The Yukawa coupling y and the mass m_{ψ^0} are dependent on the above free parameters of the model. On solving Eqs. (5.2.7) and (5.2.8) we get the expressions of these dependent parameters of our model as follows

$$y = \frac{(m_\psi - m_\chi) s_{2\alpha}}{\sqrt{2} v}, \quad (5.2.9)$$

$$\begin{aligned} m_{\psi^1} &= m_\psi c_\alpha^2 + m_\chi s_\alpha^2 = m_{\psi^0}, \\ m_\xi &= m_\psi s_\alpha^2 + m_\chi c_\alpha^2, \end{aligned} \quad (5.2.10)$$

in terms of free parameters s_α , m_ψ and m_χ .

As a result of the new interactions from the dark sector in the Lagrangian (5.2.1), we get the new gauge mediated Feynman vertices as follows

$$\begin{aligned} A_\mu \psi^+ \psi^- &: -ie\gamma_\mu, & W_\mu^+ \chi^- \bar{\psi}^0 &: -\frac{ies_\alpha}{\sqrt{2}s_W} \gamma_\mu, \\ A_\mu \chi^+ \chi^- &: -ie\gamma_\mu, & W_\mu^+ \psi^- \bar{\psi}^0 &: \frac{iec_\alpha}{\sqrt{2}s_W} \gamma_\mu, \\ Z_\mu \psi^0 \bar{\psi}^0 &: \frac{ie}{s_{2W}} \gamma_\mu, & Z_\mu \psi^+ \psi^- &: -\frac{ie(s_\alpha^2 - c_{2W})}{s_{2W}} \gamma_\mu, \\ Z_\mu \psi^+ \chi^- &: \frac{iec_\alpha s_\alpha}{s_{2W}} \gamma_\mu, & Z_\mu \chi^+ \chi^- &: \frac{ie(c_\alpha^2 - c_{2W})}{s_{2W}} \gamma_\mu. \end{aligned} \quad (5.2.11)$$

And vertices coming from the new Yukawa interactions that will play a significant role in our analysis are

$$\begin{aligned} \phi \tau^+ \psi^- &: -i(c_\alpha y_\tau^D P_R + s_\alpha y_\tau^S P_L), & h \psi^+ \psi^- &: -\frac{iy s_{2\alpha}}{\sqrt{2}}, \\ \phi \tau^+ \chi^- &: -i(-s_\alpha y_\tau^D P_R + c_\alpha y_\tau^S P_L), & h \psi^+ \chi^- &: -\frac{iy c_{2\alpha}}{\sqrt{2}}, \\ \phi \nu_\tau \bar{\psi}^0 &: -iy_\tau^D P_L, & h \chi^+ \chi^- &: \frac{iy s_{2\alpha}}{\sqrt{2}}. \end{aligned} \quad (5.2.12)$$

Apart from the those given in Eqs. (5.2.2), (5.2.7) and (5.2.8), we have only two more free parameters in our model, namely, the third generation Yukawa couplings $y_\tau^{D,S}$. We put the Yukawa couplings of light leptons at $y_e^{D,S} \sim y_\mu^{D,S} \lesssim 10^{-9}$ to conform the muon $g - 2$ measurements [225]. All the Yukawa couplings are to remain in the perturbative regime such that $y_\tau^{D,S}, y \leq 4\pi$. A conservative choice of $\lambda_{h\phi} \lesssim 10^{-4}$ is in place all through our analysis to keep the bounds from the direct detection searches and the invisible decay measurements at bay.

5.3 Relic density analysis

Due to the presence of more than one dark sector particles, the DM number changing processes, in this model, are three-fold: (i) pair annihilation ($\phi\phi \rightarrow \text{SM SM}$), (ii) coannihilation ($\phi\psi^{\pm 0} \rightarrow \text{SM SM}$), and (iii) mediator annihilation ($\psi^{\pm 0}\psi^{\mp 0} \rightarrow \text{SM SM}$). Keeping in mind the assumption of thermal freeze-out, that the dark sector particles are in equilibrium with the thermal bath in the early Universe, whereas, in chemical equilibrium with each other, one can write the Boltzmann equation describing the number density, n of the DM as follows [148].

$$\frac{dn}{dt} = -3Hn - \langle \sigma_{\text{eff}} v \rangle (n^2 - n_{\text{eq}}^2). \quad (5.3.1)$$

The effective velocity-averaged annihilation cross section, $\langle \sigma_{\text{eff}} v \rangle$ can be written as

$$\langle \sigma_{\text{eff}} v \rangle = \left(\sum_i \bar{g}_i \right)^{-2} \frac{1}{2} \sum_{i,j} \bar{g}_i \bar{g}_j \langle \sigma_{ij \rightarrow \text{SM SM}} \rangle, \quad (5.3.2)$$

where the indices i, j denote any of the dark sector particles ϕ , $\psi^{\pm 0}$ or $\chi^{\pm}$, and

$$\begin{aligned}\bar{g}_i &= g_i (1 + \delta m_i/m_\phi)^{3/2} \exp[-x \delta m_i/m_\phi], \\ \delta m_i &= (m_i - m_\phi), \quad x = m_\phi/T, \\ g_\phi &= 1, \quad g_{\psi^\pm} = g_{\chi^\pm} = g_{\psi^0} = 2.\end{aligned}\tag{5.3.3}$$

For our analyses, we have implemented the Lagrangian (5.2.1) along with all the above relations (Eqs. (5.2.2), (5.2.5) and (5.2.7) to (5.2.10)) in FEYNRULES [23]. Using the resulting model file, we carried out all the following DM analysis with the help of MICROMEGAS [50].

As discussed in the previous section, mixing affects only the charged dark fermions. Therefore $\langle \sigma_{\text{eff}} v \rangle$ will be sensitive to mixing for the annihilation channels which involve these dark fermions in the initial state and/or the propagator. One such possibility is the pair annihilation of ϕ , where the charged dark fermions appear in the t -channel propagator (Fig. 5.3.1).

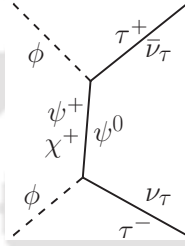


Figure 5.3.1: Pair annihilation possibilities

But the mixing effects are most prominent if one considers the coannihilation channels. This is because in some of the coannihilation diagrams (Figs. 5.3.2a) and b)), in addition to the propagators, one of the initial state particles are directly affected by mixing. Therefore, the coannihilation channels of ϕ with ψ^+ and χ^+ amply show the mixing effects.¹ On the other hand, ψ^0 not being a mixed state, $\phi\psi^0$ coannihilation channel (Fig. 5.3.2c)) is much less affected by mixing. As we are interested to study the mixing effects on the phenomenology, we will mostly concentrate on the coannihilation channels for our DM analysis. This implies that the mass splitting between DM and the dark fermions should be small throughout the study ($\lesssim 30$ GeV) and the dark sector-SM coupling is not very large ($\lesssim 1$).

For fixed mass splittings between the dark sector particles, the DM-SM Yukawa coupling vs DM mass correlation is discussed amply in the literature [57, 179]. In Ref. [91], the transition between the three kinds (pair, co- and mediator annihilation) of DM annihilation is depicted for a similar model with only a lepton doublet as the coannihilating partner of the scalar DM. In this work, this dynamics becomes more interesting in the presence of the charged singlet and doublet dark fermion mixing. In Fig. 5.3.3, y_τ^D is plotted against m_ϕ for fixed δm and y_τ^S . In the absence of mixing, depicted by the red line in Fig. 5.3.3a), the scenario is identical to Ref. [91, Fig. 4]. In this case, apart from the subdominant pair annihilation, only $\phi\psi^+$ coannihilation contributes to relic density, $\phi\chi^+$ channel being redundant due to $y_\tau^S=0$. However, as mixing increases, both $\phi\psi^+$ and $\phi\chi^+$ coannihilation channels contribute. For our choice of parameters, $\delta m_\chi \leq \delta m_\psi$, hence $\phi\chi^+$ coannihilation (Fig. 5.3.2a)) is more dominant than the $\phi\psi^+$ counterpart for a non-zero mixing. This is because both $\phi\psi^+$ and $\phi\chi^+$ coannihilations are predominantly controlled by the channels which have W^+ in the final state. Now, for $\phi\psi^+$ coannihilation case, these channel exists even for $s_\alpha = 0$, whereas, for $\phi\chi^+$ case, this channel is realised largely through mixing. Therefore, the mixing effect is more prominent in $\phi\chi^+$ coannihilation scenario over the $\phi\psi^+$ case.

With y_τ^S set to zero, as one goes for larger mixing angle, $\phi\chi^+\tau^-$ coupling (Eq. (5.2.12)) increases. This, along with the fact that both dark fermions now can non-trivially coannihilate, the total DM annihilation cross section effectively increases. Hence, for a fixed m_ϕ , larger mixing corresponds to smaller coupling to be relic density allowed. Therefore, y_τ^D gradually decreases with increase in s_α .

¹Throughout this chapter, we denote the charged dark leptons by + sign. However, All the arguments are equally applicable for negatively charged particles as well.

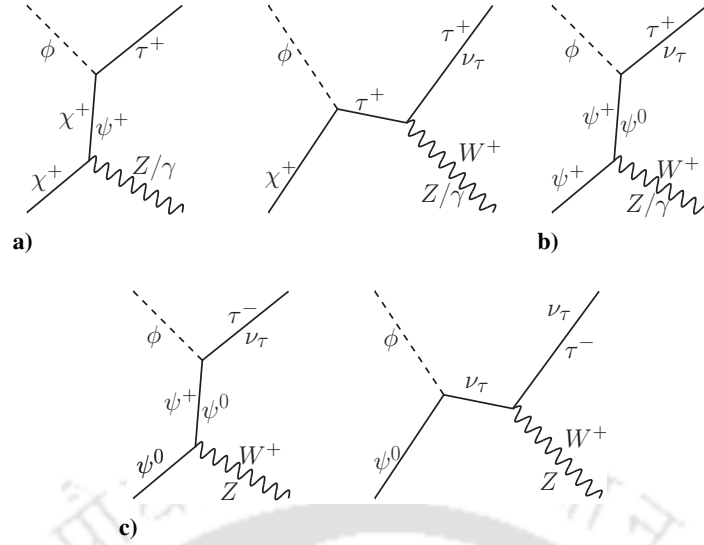


Figure 5.3.2: Coannihilation channels that are non-trivially dependent on mixing. Figs. 5.3.2a) and b) depict DM-charged fermion coannihilation whereas Fig. 5.3.2c) represents DM-neutral fermion interaction.

The scenario changes substantially if the same correlation is drawn with a fixed non-zero y_τ^S . In Fig. 5.3.3b), we assign a large value to y_τ^S . For such a large coupling, pair annihilation has a substantial contribution to the relic density. This is due to the fact that in this model, in the expression for annihilation cross section, the functional dependence on the DM-SM Yukawa coupling (λ , let us suppose) is λ^4 in pair annihilation, while for coannihilation, it is λ^2 . This is also in agreement with [91, p.5, p.6] where we showed that for a fixed δm and variable m_ϕ , coannihilation is dominant for small m_ϕ and at large DM mass, pair annihilation takes over. In this plot, as y_τ^S is already large, y_τ^D is not required to be as high as that in Fig. 5.3.3a). Now, for small m_ϕ values, $\phi\chi^+$ has the dominant contribution in relic density. $\phi\chi^+\tau^-$ coupling (Eq. (5.2.12)) is the largest for $s_\alpha = 0$ and due to this, in order to achieve right relic, the required value of y_τ^D is the smallest, as we see in this plot for the red band around $m_\phi \leq 180$ GeV. As the mixing increases, the above coupling decreases, therefore larger y_τ^D is required. However, as m_ϕ increases further, pair annihilation takes over and the mixing effect is lost in the “tail” region of the correlation around $m_\phi > 500$ GeV. As discussed above, this is because mixing does not play a significant role in the pair annihilation channels.

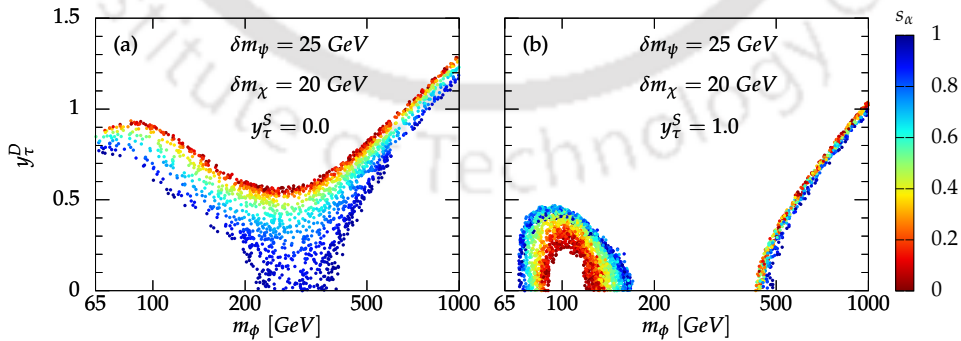


Figure 5.3.3: y_τ^D vs m_ϕ correlation for relic density allowed points. In Fig. 5.3.3a), $\phi\psi^+$ coannihilation survives for zero mixing and $\phi\chi^+$ coannihilation dominates for non-zero values of s_α . In Fig. 5.3.3b), due to large y_τ^S , pair annihilation contribution becomes prominent for the large m_ϕ .

In Fig. 5.3.3, we considered constant mass splittings between the dark sector particles. Now let us discuss how different values of these mass splittings affect the DM dynamics through mixing. In Fig. 5.3.4, y_τ^D and y_τ^S correlation is plotted as a function of s_α and for different δm 's. Similar to Fig. 5.3.3, here also we have chosen $\delta m_\chi \leq \delta m_\psi$ to facilitate the $\phi\chi^+$ coannihilation. As already dis-

cussed, large mass splitting between DM and dark fermions will suppress the coannihilation contribution towards relic density and larger coupling is required to compensate for that deficit. Hence, depending on the value of m_χ and m_ψ , the shift of the relic allowed points along y_τ^S axis is straightforward. Rather, the distribution of the points along the y_τ^D axis is an interesting feature.

For low mixing, the dominant $\phi\chi^+$ coannihilation channel is $\phi\chi^+ \rightarrow \tau^+ Z(\gamma)$ (Fig. 5.3.2a)) and its cross section is almost proportional to $(y_\tau^S)^2$. This satisfies relic density even for $y_\tau^D = 0$ as one can see from the red band in the subfigures of Fig. 5.3.4. As the mixing increases, initially, there is a contest between the two terms in $\phi\tau^+\chi^-$ coupling Eq. (5.2.12). For low y_τ^D , as mixing increases, to compensate for the overall reduction of the coupling, we see a slight increase in y_τ^S . But as y_τ^D increases further, the second term in the expression for the coupling becomes gradually negligible, making the coupling almost proportional to $s_\alpha y_\tau^D$ for large mixing. Therefore we observe that to keep within the relic bounds, for large mixing and large y_τ^D , y_τ^D decreases as s_α gradually increases. It is clear in all the plots except Fig. 5.3.4c) that the green and blue points shift to the left along y_τ^D axis as s_α gradually increases.

Now, we observed that in Fig. 5.3.4a), where $m_\psi - m_\chi = 5$ GeV, $\phi\chi^+$ coannihilation is the dominant channel. But as this splitting decreases to 2 GeV in Fig. 5.3.4b), $\phi\psi^0$ channel also contributes substantially. This is because of the coannihilation channels having W boson in the final state, which are available for $\phi\psi^0$ coannihilation even without mixing but accessible largely through mixing for $\phi\chi^+$ coannihilation. The presence of these channels suppresses the mixing dependence in the y_τ^D vs y_τ^S correlation, causing reduction in the spread of multicoloured points in Fig. 5.3.4b) compared to Fig. 5.3.4a). This spread becomes zero in Fig. 5.3.4c), where m_ψ and m_χ become almost degenerate. As already argued, here $\phi\psi^0$ coannihilation is the dominant channel in this case with $W^\pm$ in the final state and this makes the correlation completely independent of mixing.

For Fig. 5.3.4d) - f), δm_ψ is large so that $\phi\psi^+$ coannihilation contribution is very small. This establishes $\phi\chi^+ \rightarrow \text{SM SM}$ as the only active coannihilation channel. In absence of $\phi\psi^+$ (hence also $\phi\psi^0$) coannihilation and for low mixing, the only surviving $\phi\chi^+$ coannihilation in $\langle\sigma_{\text{eff}} v\rangle$ has little or no dependence on y_τ^D . The points corresponding to very low mixing, amply show this effect as the red band gradually flattens out with increasing δm_ψ as we go from Fig. 5.3.4d) - f). Ultimately in Fig. 5.3.4f), the red band widens and becomes parallel to y_τ^D axis and the blue band corresponding to large mixing becomes the thinnest.

Finally, let us see the effect of variable δm 's instead of discrete intervals discussed so far. The two plots in Fig. 5.3.5 show two such correlations of the δm 's as a measure of mixing. Fig. 5.3.5a) focusses on the asymmetry of δm_χ vs δm_ψ correlation as a measure of mixing. Mixing plays a vital role here in the distinction between the dark fermions, unlike Ref [91], where such a discrimination was not possible. The two δm 's vary over the same range and the Yukawa couplings y_τ^D and y_τ^S are fixed at the same value. This implies that apart from the mixing parameter, χ^+ and ψ^+ stand on equal footing in the context of coannihilation possibilities. However, on fixing different values for the mixing, we see that the plot shows a clear distinction between the two dark leptons of different isospins. In absence of mixing, i.e., for the blue points, the relic density allowed region along δm_χ axis broadens with the increase of s_α whereas along δm_ψ it narrows down. This is attributed to the fact that for $s_\alpha = 0$, there are more diagrams in $\phi\psi^+$ coannihilation channel than the $\phi\chi^+$ possibility. These extra diagrams (Fig. 5.3.2b)) arise due to Gauge couplings, which are exclusive to $\phi\psi^+ \rightarrow \text{SM SM}$ coannihilation because in absence of mixing, ψ^+ is purely part of the $SU(2)_L$ doublet and χ^+ is purely a singlet. It is this extra diagram that makes $\phi\psi^+$ coannihilation stronger than the $\phi\chi^+$ counterpart for $s_\alpha = 0$. This can also be verified from the distribution of the blue points, where we see that points at $\delta m_\chi = 30$ GeV correspond to δm_ψ up to 15 GeV, whereas, points at $\delta m_\psi = 30$ GeV corresponds to δm_χ only up to 12 GeV.

We have explicitly checked the above for a few benchmark points and observed that for $s_\alpha = 0$, $\phi\psi^+ \rightarrow W^+\nu_\tau$ is stronger than any of the $\phi\chi^+ \rightarrow \text{SM SM}$ channels if the two Yukawa couplings are equal. For $s_\alpha = 1/\sqrt{2}$, on the other hand, $\phi\psi^+$ and $\phi\chi^+$ coannihilation contribution becomes equal even for these extra Gauge channels, so the red points show a symmetric distribution along both the axes. For the green points, i.e., for $s_\alpha = 1$, we simply see the opposite of the $s_\alpha = 0$ case, because

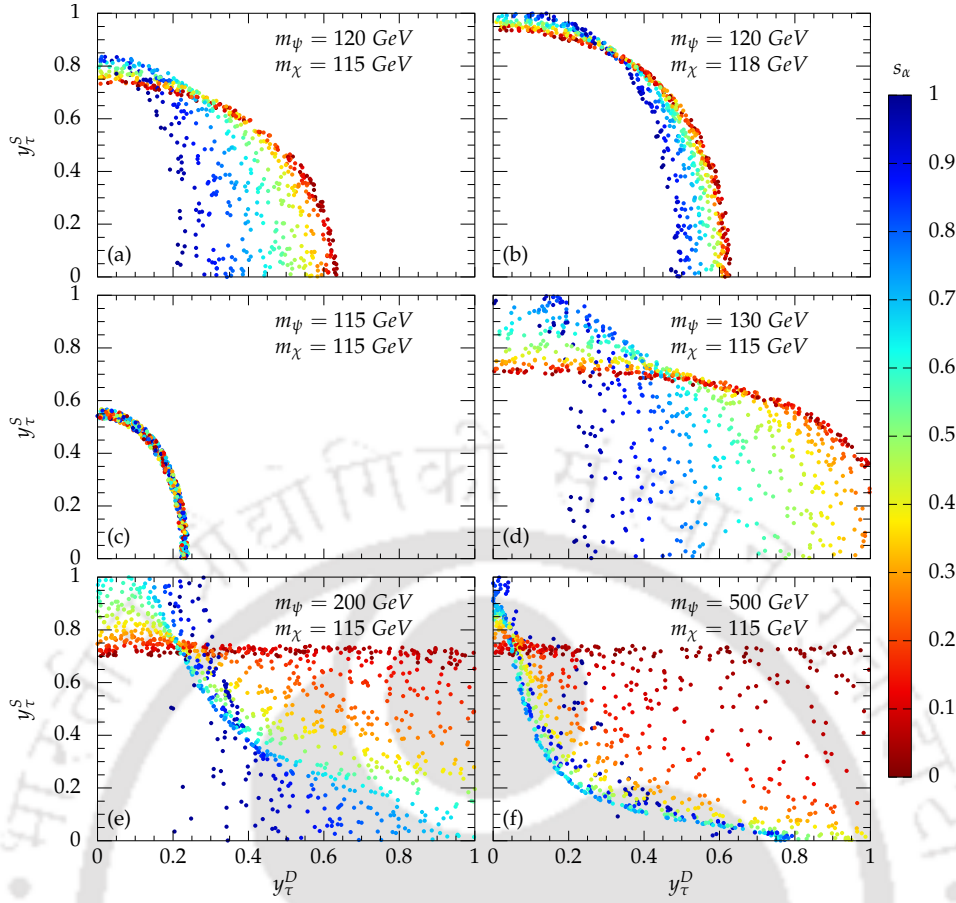


Figure 5.3.4: y_τ^D vs y_τ^S correlation for relic density allowed points. For fixed DM mass, the variation of the charged dark fermion masses and mixing angle substantially changes the correlation. m_ϕ is fixed at 100 GeV.

$W^+ \psi^- \psi^0|_{s_\alpha=0} = W^+ \chi^- \psi^0|_{s_\alpha=1}$. This also agreed with our observation of benchmark point results, that the $\phi \chi^+$ coannihilation is more dominant than $\phi \psi^+$ counterpart for $s_\alpha = 1$. To sum up, we can conclude that the relic density allowed region for the pure leptonic eigenstates widens along δm_χ axis, whereas the mixed states tend to widen gradually along δm_ψ axis.

Fig. 5.3.5b) on the other hand, discusses $\delta m_{\psi 0}$ vs δm_ψ correlation as a function of the mixing parameter. The DM mass varies in the same range as Fig. 5.3.5a), and the Yukawa couplings are fixed at $y_\tau^D = y_\tau^S = 0.5$.

For $s_\alpha = 0$, $m_{\psi 0} = m_\psi$, as obvious from Eq. (5.2.8). Therefore, the nature of the correlation will be a straight line with a 45° slope, which is exactly what the orange line represents. However, for the non-zero s_α values, we see distribution on both sides of the red line. As we can see, the spread of the points varies with s_α .

It is to be noted that if the relic density constraint is not imposed, then for a fixed s_α , the points would be distributed symmetrically around the orange line, maintaining it as the line of symmetry. This is obvious from the model description arguments of Section 5.2. However, on adding the constraint, one can see from the plot that this symmetry is lost. For large δm 's, we observe that the relic allowed region shifts slightly to the left of the orange line for all non-zero s_α values.

We numerically checked that for fixed Yukawa couplings, as mixing increases in the range $0 \leq s_\alpha \leq 1/\sqrt{2}$, the contribution of $\phi \psi^+ \rightarrow \text{SM SM}$ channels rapidly decreases, but $\phi \psi^0 \rightarrow \text{SM SM}$ remains much less affected, except a very small increase. Our relic density calculation for a few benchmark points within the allowed region confirmed that $\phi \psi^0 \rightarrow W^+ \tau^-$ channel is the dominant channel in $\langle \sigma_{\text{eff}} v \rangle$ for this s_α window.

This explains the shift at large δm for the non-zero s_α values. As $\phi \psi^+$ contribution decreases, to

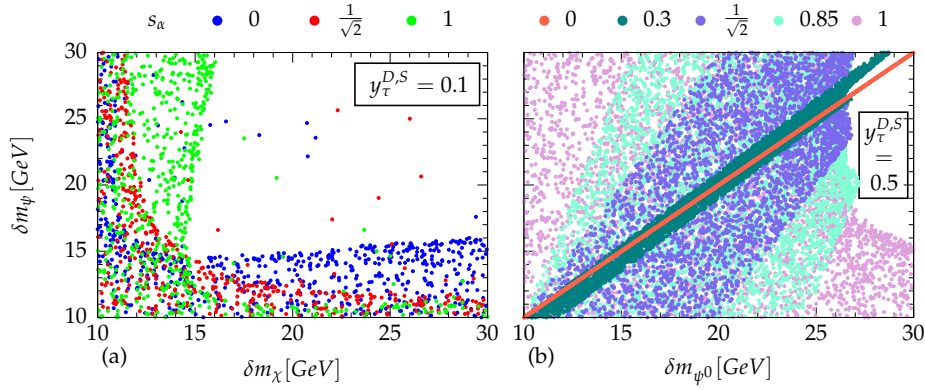


Figure 5.3.5: Correlations δm_χ vs δm_ψ in Fig. 5.3.5a) and δm_{ψ^0} vs δm_ψ in Fig. 5.3.5b) for relic density allowed points as a function of the mixing angle. Since all other parameters are the same for the dark leptons discussed in each plot, it is only up to the mixing parameter to determine the dominant relic density contribution between them. The DM mass varies in the range $65 \text{ GeV} \leq m_\phi \leq 1 \text{ TeV}$.

maintain relic abundance in the observed range, one needs stronger $\phi\psi^0$ coannihilation. This is why the region with a very large δm_{ψ^0} remains out of the relic allowed regime, the annihilation cross section being insufficient, leading to overabundance. It is also justified that we see the upper limit of relic allowed δm_0 decreasing gradually as s_α increases from 0 to $1/\sqrt{2}$.

On the other hand, we observed that $\phi\chi^+ \rightarrow \text{SM SM}$ contribution, which was negligible so far, becomes substantial for large mixing, i.e., in the range $1/\sqrt{2} \leq s_\alpha \leq 1$. As already explained, $\phi\psi^0$ coannihilation remains mostly unaffected by mixing. Therefore, as $\phi\chi^+$ contribution gradually increases with mixing in the above range, it also relaxes the exclusion limit for large δm_{ψ^0} , which can be seen for the cyan and pink points. Ultimately, for $s_\alpha = 1$, the spread of the relic allowed region is maximum. This is because the relic density is almost entirely dominated by $\phi\chi^+$ coannihilation here, so even large δm_ψ and δm_{ψ^0} values are allowed except for the white excluded region on the right, where both the δm 's are very large, leading to insufficient $\langle\sigma_{\text{eff}} v\rangle$.

The above features are model-independent and can be generalised for any singlet DM coannihilation scenario that involves a singlet as well a doublet coannihilating partner.

5.4 Experimental signatures

5.4.1 DM signatures at the LHC

The problem of finding signatures of DM in a collider environment is a very challenging task. Disentangle the DM signals from a multitude of invisible particles is notoriously hard. Several search strategies are there for this task which addresses the problem from the perspective of the hadron colliders (LHC and future hadron colliders) as well as lepton colliders which are for the very purpose where hadron collider has limitations. In hadron colliders, we do not have access to the longitudinal component of the missing momenta. Hence the task of finding the signature of DM is more difficult as we have to depend entirely on the observables constructed from the transverse momentum components. One can separate *multilepton + missing energy* signals from the backgrounds even for a small signal cross section through advanced techniques of Multivariate Analysis [91]. Here we look for specific distributions through which we can decipher the DM signals in a collider environment. We discussed previously that the presence of a doublet along with a singlet dark fermions and the mixing between them plays a significant role in the phenomenology of our model. Here we point to the ways to find out those signatures in collider environments that will highlight this feature of the model and try to distinguish its effect from the distributions. The peaks and end-points of a kinematic distribution can be associated with the masses of the mediating particles. The kinematic distributions of transverse momentum, p_T , transverse mass, m_T and invariant mass are a few very significant distributions to study.

In the following we will discuss the relevant distributions in a hadron collider environment like LHC. To perform the analysis for the LHC at the CM energy $\sqrt{s} = 13$ TeV, we proceeded as follows: (1) FEYNRULES [23] has been used to generate model files. (2) Events have been generated using MADGRAPH5 [26] and showered with PYTHIA 8 [215]. We used the dynamic factorisation and renormalisation scale for the signal as well as the background events. (3) The detector simulation has been performed with the help of DELPHES [134]. (4) The distributions were drawn with the help of MADANALYSIS 5 [109].

In the following analysis of *multilepton + missing energy* signals at the LHC, our focus is particularly on the distributions of the τ 's. As we discussed in Section 5.2, the couplings of the first two generations of the leptons are very constrained by the muon $g - 2$ measurements, whereas the τ couplings are unbounded. Although phenomenologically interesting in our case, the study of τ poses many challenges as well. As it is heavier than most of the light quark mesons, τ can decay hadronically, unlike other leptons. Moreover, the leptonic decays of τ are difficult to distinguish from prompt leptons in an $\ell + E_T^{miss}$ final state. This makes the hadronically decaying τ 's suitable for the collider signatures. We reconstructed the τ -tagged jets from DELPHES [134] with the help of FASTJET [83] using the anti- k_T algorithm, with the separation ΔR of two adjacent τ -jets to be 0.4. The distance between two objects i and k is defined as $\Delta R_{ik} = \sqrt{(\eta_i - \eta_k)^2 + (\phi_i - \phi_k)^2}$, where η_i and ϕ_i are the rapidity and azimuthal angle of the object i , respectively. In Table 5.4.1 below, we give the selection criteria, which has been used throughout the following analysis. The τ -tagging efficiency is assumed to be flat and taken to be 60%, with fake rates from light jets to the τ -jets of 1%. The statistical significance is given by $S = \sqrt{2[(S+B)\ln(1+S/B) - S]}$, where S and B respectively are the number of expected signal and background events at a particular integrated luminosity $\mathcal{L}$.

Selection parameter	$p_T^{\ell(j)}$	$ \eta_{\ell(j)} $	ΔR_{ik}
Cut value	10 (20) GeV	5	0.4

Table 5.4.1: The selection criteria used throughout the analysis, where $\ell = e, \mu$ and $j = u, d, c, s$ and g . The distance between two objects i and k is defined as $\Delta R_{ik} = \sqrt{(\eta_i - \eta_k)^2 + (\phi_i - \phi_k)^2}$, where η_i and ϕ_i are the rapidity and azimuthal angle of the object i , respectively.

For the effect of mixing angle α in the collider signatures, two kinds of processes will be useful (1) the processes that are predominantly W boson mediated, and (2) those that are also predominantly Z boson mediated.

$3\tau + E_T^{miss}$ channel

From Eq. (5.2.11), we can see that the couplings of W boson with the dark fermions are functions of mixing angle α . Without the mixing, i.e., when $s_\alpha = 0$, ψ^+ is exclusively a doublet and χ^+ purely a singlet. This fact immediately implies that when $W^+ \rightarrow \psi^0 \chi^+$ channels are off, the resulting final states are a consequence of pure doublet contribution. Hence, by tuning the mixing, one can control the % of singlet contribution in the channels. The channels we can look into are as follows:

$\tau \nu 2\phi$: s -channel processes via $W^+ \rightarrow \psi^0 \psi^+(\chi^+)$ followed by ψ^0 decaying into $\nu \phi$, which remains invisible, and $\psi^+(\chi^+) \rightarrow \phi \tau^+$. Since the only visible final state in this channel is a single τ , and the missing energy can come from both ψ^0 and $\psi^+(\chi^+)$, it is difficult to conclude anything about the DM signature.

$3\tau \nu 2\phi$: Here, the signal processes can proceed through the following different modes:

- (i) W -mediated s -channel processes through $q\bar{q}' \rightarrow \psi^0 \psi^+(\chi^+)$ channel; and
- (ii) W -mediated s -channel as well as quark-mediated t -channel processes through $q\bar{q}' \rightarrow W Z(\gamma)$ and $q\bar{q}' \rightarrow W h$ channels.

We encounter the latter case in the context of the *Unitarity problems* of the gauge bosons and is very similar to the $f\bar{f}' \rightarrow W Z(\gamma)$ and $f\bar{f}' \rightarrow W h$. See for example Refs. [101, 117, 142] for some recent papers in this context and the references therein. The point is that as the CM energy increases, the cross section of these channels decreases, which we can explain from the *Equivalence Theorem* of the gauge bosons. We see a similar trend in our case also. As a result, the dominant channels of $3\tau\nu 2\phi$ process become mostly W -mediated s -channels mentioned in point (i) above. The diagrams mentioned in point (ii) are suppressed to point (i). See Fig. 5.4.1 for the complete diagrams with their decay channels.

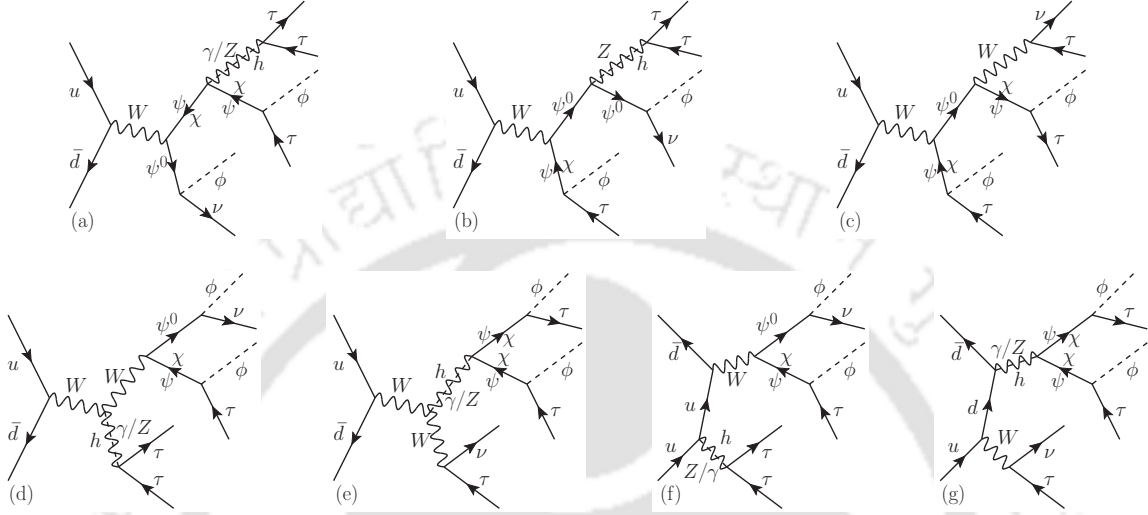


Figure 5.4.1: Feynman diagrams contributing to the $3\tau\nu 2\phi$ channel at the LHC. (i) The upper panel shows the W -mediated s -channel processes through $q\bar{q}' \rightarrow \psi^0 \psi^+(\chi^+)$ channel; and (ii) the lower panel, W -mediated s -channel as well as quark-mediated t -channel processes through $q\bar{q}' \rightarrow W Z(\gamma)$ and $q\bar{q}' \rightarrow W h$ channels. The diagrams of the lower panel are seen to be suppressed, which makes the $3\tau\nu 2\phi$ channel predominantly W -mediated.

In the CM frame, the two particles are emitted with equal and opposite transverse momentum (p_T), whose magnitude is, say p_T^{CM} . As we can see from Fig. 5.4.1, one of these two particles decays into two-body final states, whereas the other one decays into four-body final states, resulting into 3τ 's and missing energy signal. If we distinguish these three τ leptons according to the decreasing order of their p_T , the leading- p_T τ will mostly come from the two-body decays, i.e., Fig. 5.4.1b) and (c). On the other hand, the two subleading τ 's will probably come from the four-body decays mentioned above. As the third τ is the least energetic one, it will be difficult to explain the kinematics of the third τ with multiple accompanying invisible particles coming from such four-body decays. Moreover, the statistics of the third τ also comes out to be less than the other two τ 's². Hence we will confine our discussions within the two leading- p_T τ 's and their observables. It is obvious from Fig. 5.4.1 that, the two-body visible channel is $\psi^+(\chi^+) \rightarrow \phi \tau^+$ (Fig. 5.4.1b) and (c)), whereas the invisible channel is $\psi^0 \rightarrow \nu \phi$ (Fig. 5.4.1a)). In Fig. 5.4.2, we show some of the relevant distributions of the 3τ channel for a set of benchmark points (see Table 5.4.2) which satisfy the required relic density. Here the relic density is satisfied through pair annihilation channels as we concluded for the doublet case [91].

From the distributions of Fig. 5.4.2, we can deduce the following inferences:

(i) We see sharp ‘‘Jacobian peaks’’ in the p_T distributions for leading and subleading- p_T τ 's for BPs 5-7. These ‘‘Jacobian peaks’’ appear at $m/2$ in p_T -distributions and at m in missing transverse mass distributions. Here m is the mass of the parent particle which decays to τ . From the peaks of Fig. 5.4.2a) and (b) we can infer that τ_1 is coming from χ and τ_2 from W bosons for BP7. We can conclude that both χ and W boson are on-shell for this benchmark point. For BPs 5-6, the peaks are slightly towards

²This assertion is purely based on our Monte Carlo data sample. We showed, in Fig. 5.4.2, some representative distributions which can be useful in depicting the effects of mixing. To make a conclusive statement about the features of the third τ in this process, we need a very large data set of Monte Carlo events. This difficulty is more so because of τ detection efficiency at the LHC.

	m_ϕ	m_{ψ^0}	$m_{\psi^\pm}$	m_χ	y_τ^D	y_τ^S	y	s_α	cross section [fb]
	[GeV]								
BP1	100	348	350	180	0.4	2.05	0.097	0.1	1.2
BP2	120	249	250	150	2.15	0.0	0.057	0.1	0.15
BP3	180	358	380	295	0.34	2.0	0.2	0.5	0.10
BP4	180	358	380	295	0.72	0.85	0.2	0.5	0.058
BP5	80	250	350	150	1.87	0.0	0.57	0.71	2.1
BP6	80	250	350	150	0.55	2.5	0.57	0.71	12.2
BP7	88	1406	1500	110	2.8	0.0	2.0	0.26	0.049

Table 5.4.2: Relic density allowed benchmark points for analysis of $3\tau\nu 2\phi$ channels. $y_\ell^D = y_\ell^S = 10^{-9}$. The shaded columns are for the dependent model parameters.

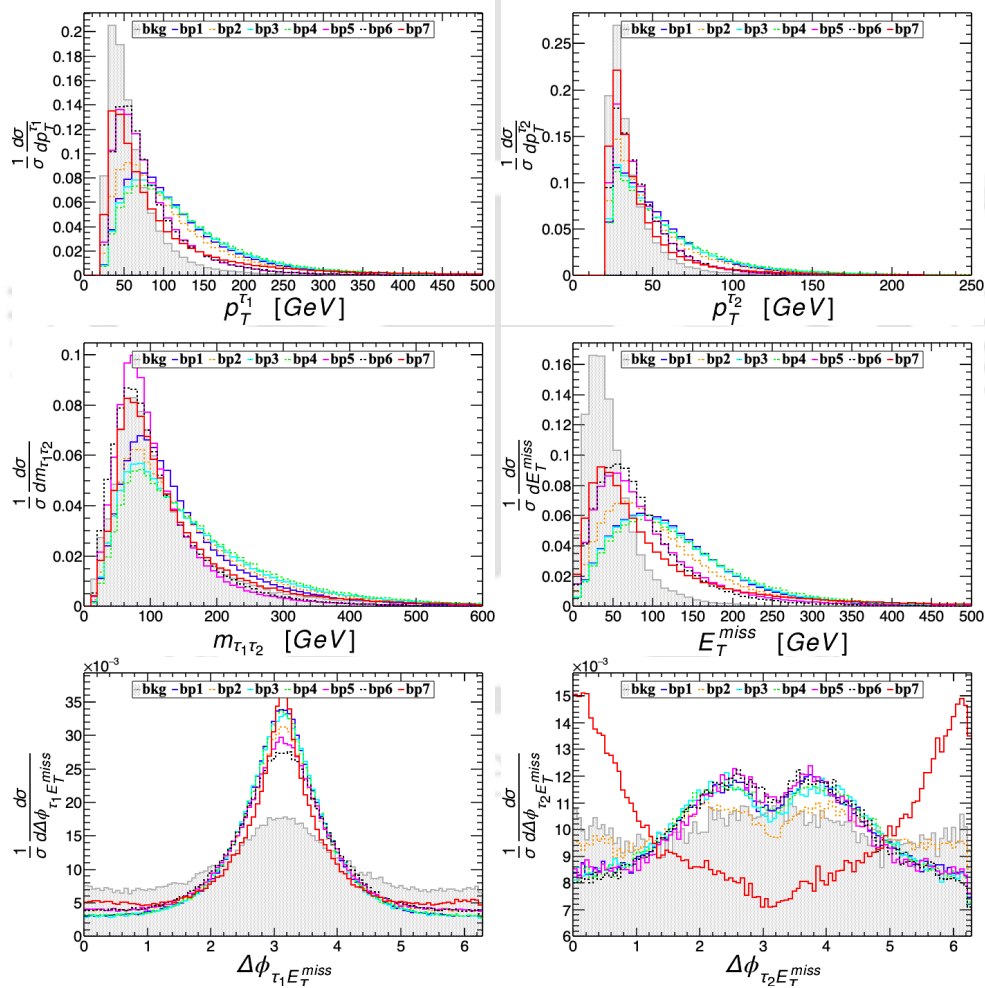


Figure 5.4.2: Significant distributions of the $3\tau\nu 2\phi$ channel. Gradual shift of the peaks in the top and middle panel can be attributed to the change in the mixing parameter. $\Delta\phi$ distributions shows direction of the production of missing energy signal w.r.t. τ 's, which are clearly distinguishable from the background.

the right. This fact signifies that here the leading and sub-leading τ 's do not entirely come from χ and W , but predominantly so. We see further confirmation for this from the distributions of the respective

transverse masses which we have not included for the sake economy of space.

(ii) The BPs 1-4 and 5-7 can be similarly classified. While the earlier set gives a relatively flatter profile, the latter shows sharp peaks in the p_T and E_T^{miss} distribution plots. From the discussion above, we can see that the reason for this is the off-shell-ness of the mediating particles. This point is also clear from the values of the parameters in Table 5.4.2.

(iii) The $p_T^{\tau_1}$ distribution profile is similar to that of E_T^{miss} , whereas the peak of $\Delta\phi_{\tau_1 E_T^{miss}}$ shows that τ_1 and E_T^{miss} are back to back in nature. We can infer that they are coming from the same parent particles and so they are equal and opposite in nature.

(iv) The invariant mass plot of τ_1 and τ_2 gives the signal that they are pair produced from the Z boson decay.

(v) The distinct features of BP 7 are evident from all the distributions. The very high values of $m_{\psi^{0\pm}}$ indicate that these particles in the intermediate states will suppress the effects of the respective diagram very much compared to χ .

$\ell \tau + E_T^{miss}$ channel

Now we shall look into the effects of mixing in the Z boson mediated signal processes. For this, we focus on the $\ell \tau \nu 2\phi$ channel. The processes in this channel can proceed through the following different modes, as can be seen from Fig. 5.4.3:

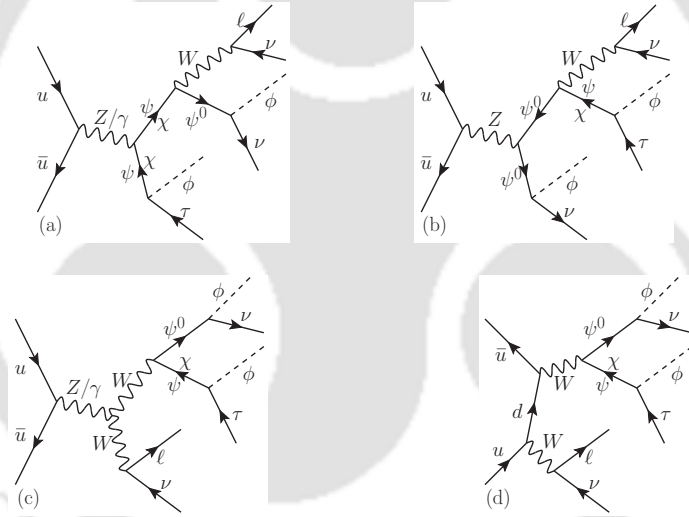


Figure 5.4.3: Feynman diagrams contributing to the $\ell \tau \nu 2\phi$ channel at the LHC. (i) The upper panel shows the Z -mediated s -channel processes through $q\bar{q} \rightarrow \psi^+ \psi^-, \chi^+ \chi^-, \psi^+ \chi^-$ and $\psi^0 \bar{\psi}^0$ channel; and (ii) the lower panel, Z -mediated s -channel as well as quark-mediated t -channel processes through $q\bar{q} \rightarrow W W$ channel. The diagrams of the lower panel are seen to be suppressed, which makes the $\ell \tau \nu 2\phi$ channel predominantly Z -mediated.

- (i) Z boson mediated s -channel processes through $q\bar{q} \rightarrow \psi^+ \psi^-, \chi^+ \chi^-, \psi^+ \chi^-$ and $\psi^0 \bar{\psi}^0$ channel; and
- (ii) Z boson mediated s -channel as well as quark-mediated t -channel processes through $q\bar{q} \rightarrow W W$ channel.

Among these, the processes in point (i) are much dominant to the processes in point (ii). This is because, as in the case for $3\tau \nu 2\phi$ channel, here also the latter case is very similar to the $f\bar{f} \rightarrow W W$ in the context of the Unitarity problems of the gauge bosons [101, 117, 142]. Hence, as the CM energy increases, the cross section of these channels decreases, that can be explained from the Equivalence Theorem of the gauge bosons. Among the remaining diagrams of point (i), the γ mediated processes are independent of the mixing parameter s_α at the production level, as can be seen from Eq. (5.2.11). As a result, they

are not interesting for our purpose, where we focus on the effects of the mixing parameter in the distributions. The $\ell \tau 2\nu 2\phi$ channel can also have contributions from gg initiated s -channel processes which are very much suppressed for the chosen benchmark points (see Table 5.4.3). The $q\bar{q}$ initiated Higgs boson channels are also suppressed due to negligible hqq couplings. So, it should be sufficient to focus only on the processes in point (i) above, while explaining the features of this channel. In Fig. 5.4.4, we show the relevant distributions of the $\ell \tau 2\nu 2\phi$ channel for the set of benchmark points in Table 5.4.3 which satisfy the required relic density.

	m_ϕ	m_{ψ^0}	$m_{\psi^\pm}$	m_χ	y_τ^D	y_τ^S	y	s_α	cross section [fb]
	[GeV]								
BP		300				0.23	0.57	$1/\sqrt{2}$	4.4
BP_5		350				0.26	0.5	0.5	1.88
BP_1	120	398	400	200	0.5	1.55	0.11	0.1	0.73
BP_05		399.5				2.2	0.0574	0.05	0.35
BP_03		399.82				2.12	0.0345	0.03	0.25
BP_0		400				1.82	0.0	0.0	0.18

Table 5.4.3: Relic density allowed benchmark points for analysis of $\ell \tau 2\nu 2\phi$ channels. $y_\ell^D = y_\ell^S = 10^{-9}$. The shaded columns are for the dependent model parameters.

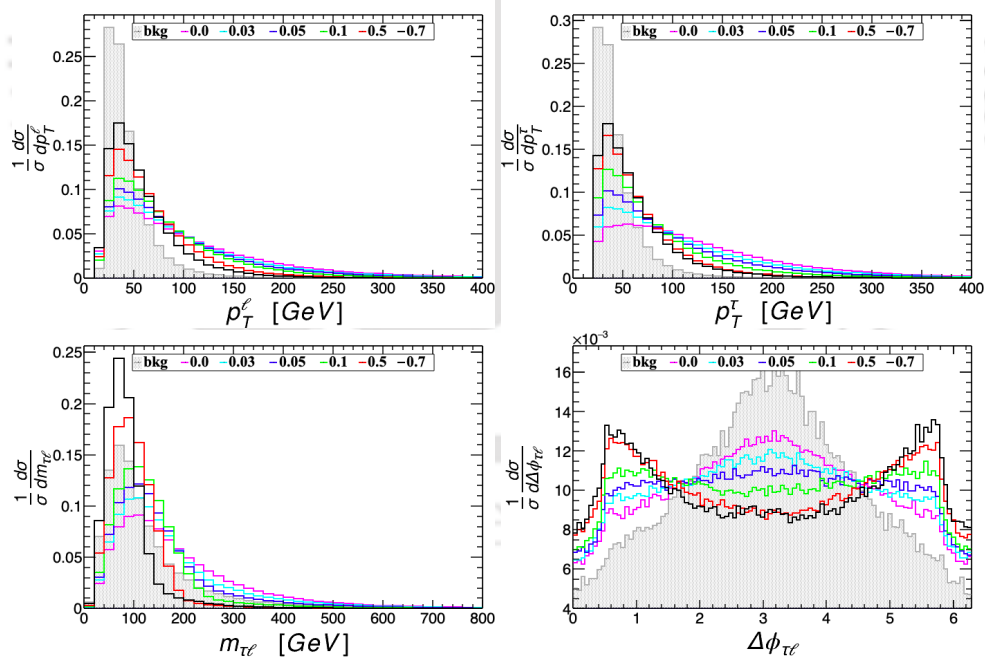


Figure 5.4.4: Some significant distributions of the $\ell \tau 2\nu 2\phi$ channel. As for the $3\tau \nu 2\phi$ case, gradual shift of the peaks in the distributions can be attributed to the change in the mixing parameter. The distributions are labelled with the respective values of s_α .

Some interesting observations from the distributions of Fig. 5.4.4, are as follows:

(i) If we closely look into the benchmark points for this channel in Table 5.4.3, they are chosen such that the first one is for $s_\alpha = 1/\sqrt{2}$ where there is equal mixing between the doublet and singlet component of the dark fermions. We gradually change the value of s_α to zero. The rest of the dependent and independent parameters are chosen just to keep the relic density within the allowed limit.

(ii) For $s_\alpha = 0$, there is no mixing, and hence ψ^+ is a pure doublet state whereas χ^+ is purely singlet. Now since we see from Eq. (5.2.11) that $W^+\chi^-\psi^0$ and $Z\psi^+\chi^-$ vertices do not exist for such a case, we have only the doublet contribution for BP_0. As a result, the distributions are independent of the values of m_χ and y_τ^S .

(iii) From the numbers given in Table 5.4.3, we see that the role of mixing dictates the dominant channel, and hence, the trends of invariant mass distribution. Thus, we conclude from these numbers that, Fig. 5.4.3b) dominates for the values of s_α around $1/\sqrt{2}$ leading to a Z -peak for $s_\alpha = 1/\sqrt{2}$, whereas, for other values of s_α Fig. 5.4.3a) dominates. This can also be seen in the distribution of $\Delta\phi_{\tau\ell}$ in Fig. 5.4.4, where we see that both τ and ℓ are more probable to be along the same direction around $s_\alpha = 1/\sqrt{2}$. $\Delta\phi_{\tau E_T^{miss}}$ further establishes this conclusion where we will find τ and E_T^{miss} going in the opposite directions. We have not kept $\Delta\phi_{\tau E_T^{miss}}$ plot here to avoid redundancy.

(iv) The p_T distributions of Fig. 5.4.4 justify the dominance of the channels of Fig. 5.4.3 as a function of the mixing angle. In both p_T^ℓ and p_T^τ distributions, it is seen that towards the low p_T region, Fig. 5.4.3b) dominated channels are more probabilistic whereas the high p_T region favours Fig. 5.4.3a) dominated channels. This could be attributed to the feature that in Fig. 5.4.3b) dominated channels, the energy share of both the light lepton and τ is less than the energy share in Fig. 5.4.3a). In Fig. 5.4.3b), both the lepton and τ are produced from the decay of a single ψ^0 whereas in Fig. 5.4.3a), they are produced from different mediators. This feature is more prominent in p_T^τ rather than p_T^ℓ . This is because in Fig. 5.4.3a), the branching of the decay chain producing τ is comparatively less than what it is in Fig. 5.4.3b), making the difference in energy share of τ between Fig. 5.4.3a) and (b) dominated channels more prominent. The lepton, on the other hand, is produced through more branching in both the figures, which makes the effect of reduced energy share less prominent as Fig. 5.4.3b) takes over Fig. 5.4.3a) with the increase of mixing angle in the chosen benchmark points.

We observe the same trend through all the distributions of Figs. 5.4.2 and 4. There is a gradual change in the peak and tail positions of the kinematic distributions with the variations of s_α . This change is more prominent for the latter case as we have chosen the benchmark points as such. Here, apart from the mixing parameter, s_α , all the other dependent and independent parameters of the model are tweaked to some extent so that they satisfy the required relic density. Despite these small tweaks, we can say from our observations that this feature of the distributions is solely dependent on s_α and not anything else. We confirmed our assertion by keeping the rest of the independent parameters same and varying only the value of s_α .

5.4.2 Indirect Search Prospects

Apart from the discovery potential at the collider discussed above, the dark sector mixing can affect other DM detection possibilities as well. For the leptophilic scalar DM discussed here, the dominant indirect detection channel is $\phi\phi \rightarrow \tau^+\tau^-$ (Fig. 5.4.5). We discussed the Fermi-LAT constraints on parameter space for the scalar DM and lepton doublet interaction in Ref. [91, p.7]. There, we showed that most of the parameter region is allowed by experimental bound, except for a small region at low DM mass. In $m_\phi \lesssim 200$ GeV, the region with large Yukawa coupling ($y_\tau \gtrsim 2.0$) is excluded by Fermi-LAT limits and this bound becomes more stringent if one considers small Δm .

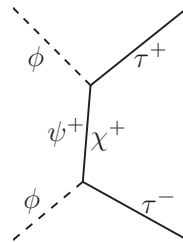


Figure 5.4.5: DM indirect search channel. Dark fermions appear in the propagator, facilitating the mixing dependence.

In this work, the velocity-averaged annihilation cross section, i.e., $\langle\sigma v\rangle_{\tau\tau}$ will show dependence on

the mixing parameter as the charged dark leptons appear in the t -channel propagator. In Table 5.4.4, we show that for fixed δm 's and Yukawa couplings, a finite mixing can bring $\langle\sigma v\rangle_{\tau\tau}$ below the Fermi-LAT bounds which would otherwise be above the limits in [91].

	m_ϕ	$m_{\psi^\pm}$	m_χ	y_τ^D	y_τ^S	s_α	$\langle\sigma v\rangle_{\tau\tau}$ [cm ³ /s]
	[GeV]						
BP1	100	600	110	2.5	0.0	0.03	8.64×10^{-28}
BP2	105	600	130	2.0	0.0	0.45	6.23×10^{-27}
BP3	125	300	140	1.75	0.03	0.25	1.39×10^{-27}
BP4	150	400	175	2.5	0.0	0.27	5.98×10^{-27}

Table 5.4.4: Relic density allowed benchmark points which are below the Fermi-LAT limit. Around $m_\phi \sim 100$ GeV, the Fermi-LAT bound is $\sim 10^{-26}$ cm³/s.

The bound on the upper limit of Yukawa coupling can be relaxed for a finite dark sector mixing. It is obvious that due to having two charged dark fermions in this work, there are two $\Delta m (= \delta m/m_\phi)$ parameters. On keeping one Δm small and setting the other at a high value, it is possible to address both relic density and indirect search constraints. It becomes very interesting to apply this in low m_ϕ and small Δm scenario, because as said previously, this region is typically above the limits in similar models.

The benchmark points in Table 5.4.4 are so chosen that for each point, DM-SM interaction takes place predominantly through the mixing. The relic density, on the other hand, is satisfied through various coannihilation and mediator annihilation channels which also have a strong mixing dependence. We observe here that for low DM mass of around 100 GeV, $\langle\sigma v\rangle_{\tau\tau}$ remains below the Fermi-LAT limit even for coupling as high as $y_\tau^D \sim 2.5$ and one of the Δm 's sufficiently small.

5.4.3 Combined parameter region

From the discussions above, the effect of dark matter mixing is evident in all the signatures of this model. In Section 5.4.1 we discussed the collider signatures for the $3\tau + E_T^{miss}$ and $\ell\tau + E_T^{miss}$ channels through different kinematic distributions, whereas Section 5.4.2 gives us the effect of mixing in the indirect search experiments. Here, in this section, we present a combined analysis where we take into account all the constraints together and see the prospects of detection of such signatures in the current and future experimental facilities.

Following the definition of the *fiducial phase space* in Ref. [1], the criteria for the same are the selection of *three τ -jets (one light charged lepton and one τ -jet) plus non-zero E_T^{miss} for $3\tau + E_T^{miss}$ ($\ell\tau + E_T^{miss}$) channel along with the cuts mentioned in Table 5.4.1. The *fiducial cross section* is defined as*

$$\sigma_{\text{fid}} = \frac{N - B}{\epsilon \cdot \mathcal{L}} \quad (5.4.1)$$

where $\mathcal{L}$ is the integrated luminosity, N is the expected number of observed events, B is the estimated number of background events and ϵ is the efficiency of the detector to separate the particular signal events from the backgrounds. B includes the background events from the final states given in Table 5.4.5. Since the SM sources of $3\tau + E_T^{miss}$ ($\ell\tau + E_T^{miss}$) channel are also taken among the background processes, the numerator of Eq. (5.4.1) gives us the excess number of events for the new physics signal. Thus, we can write, $N - B = \epsilon \cdot \sigma_{\text{fid}} \cdot \mathcal{L} = S$. We estimated ϵ in our simulation as the ratio of the number of signal events passing the selection requirements at detector level to those passing the fiducial selection criteria listed in Table 5.4.1 at the parton level. Hence, using the τ -tagging efficiency and faking rate mentioned

earlier, we calculate N , using Eq. (5.4.1), which gives us the estimated number events necessary to see the difference between the distributions of different benchmark points at a particular integrated luminosity.

	backgrounds	cross section [pb]	events	total events
$3\tau + E_T^{miss}$	$t\bar{t} \rightarrow \tau 2j \nu b\bar{b}$	66.275	397	33069
	$WW \rightarrow \tau 2j \nu$	10.073	60	
	$WZ(\gamma^*) \rightarrow \tau 2j \nu$	0.125	29908	
	$WZ(\gamma^*) \rightarrow 3\tau \nu$	4984.69	2704	
$\ell\tau + E_T^{miss}$	$t\bar{t} \rightarrow \ell\tau 2\nu b\bar{b}$	33.14	19884	327626
	$t\bar{t} \rightarrow \ell 2j \nu b\bar{b}$	199.194	3983	
	$WW \rightarrow \ell\tau 2\nu$	5.026	3015	
	$WW \rightarrow \ell 2j \nu$	30.239	604	
	$WZ(\gamma^*) \rightarrow \ell 2\tau \nu$	0.376	451	
	$WZ(\gamma^*) \rightarrow \ell 2j \nu$	14972.851	299457	
	$WZ(\gamma^*) \rightarrow 2\ell\tau \nu$	0.388	232	

Table 5.4.5: Estimate of the background events for the integrated luminosity, $\mathcal{L} = 100 \text{ fb}^{-1}$.

Taking all the considerations mentioned above, in Fig. 5.4.6, we show the variation of the statistical significance $\mathcal{S}$ as function of the mixing angle α for fixed values of other parameters, namely, $(m_\phi, m_{\psi^\pm}, m_\chi) = (120, 420, 150) \text{ GeV}$ and $(y_\tau^D, y_\tau^S) = (0.4, 0.21)$. We have chosen three different values of $\mathcal{L} = 100, 300 \text{ fb}^{-1}$ and 3 ab^{-1} for our study. We observe that $\mathcal{S}$ increases as s_α approaches $1/\sqrt{2}$. From Fig. 5.4.6, it is also evident that the $3\tau + E_T^{miss}$ channel gives us better chance of detecting the mixing signatures for a larger region of parameter space than the $\ell\tau + E_T^{miss}$ channel.

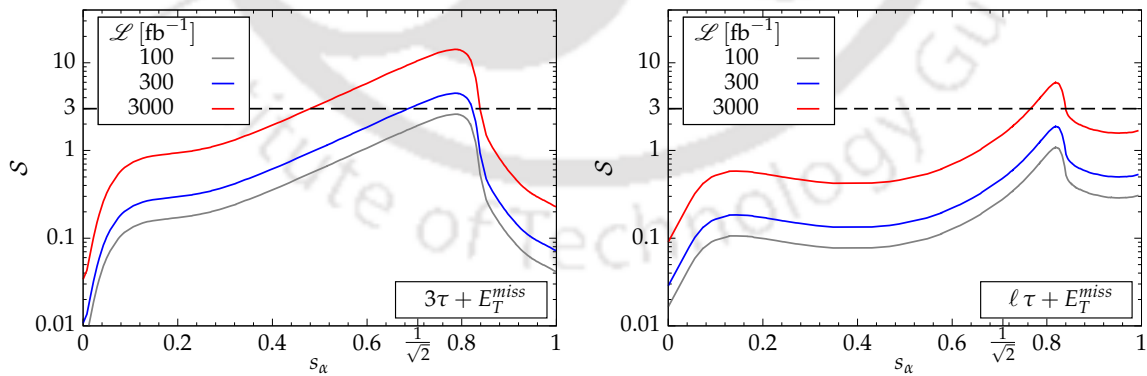


Figure 5.4.6: The variation of the statistical significance, $\mathcal{S}$ with the mixing parameter s_α for fixed values of other free parameters [$(m_\phi, m_{\psi^\pm}, m_\chi) = (120, 420, 150) \text{ GeV}$ and $(y_\tau^D, y_\tau^S) = (0.4, 0.21)$]. $\mathcal{S} = \sqrt{2}[(S+B)\ln(1+S/B) - S]$, where S and B respectively are the number of expected signal and background events at a particular integrated luminosity $\mathcal{L}$.

Therefore, we perform a combined analysis and show the s_α vs m_ϕ correlation in Fig. 5.4.7, taking into account all the experimental constraints. The shaded regions represent the allowed parameter space, with the blue, red and green shade corresponding to indirect detection, relic density and collider limits respectively. Here, we have shown the collider limits for $3\tau + E_T^{miss}$ channel at $\mathcal{L} = 3 \text{ ab}^{-1}$. This corresponds to $\mathcal{S} \geq 3$ for $0.46 \leq s_\alpha \leq 0.83$, as seen for the green shaded area. However, for smaller s_α

in the above window, one can see that very large m_ϕ is not allowed. The low value of $\mathcal{S}$ in this region is consistent with that of Fig. 5.4.6.

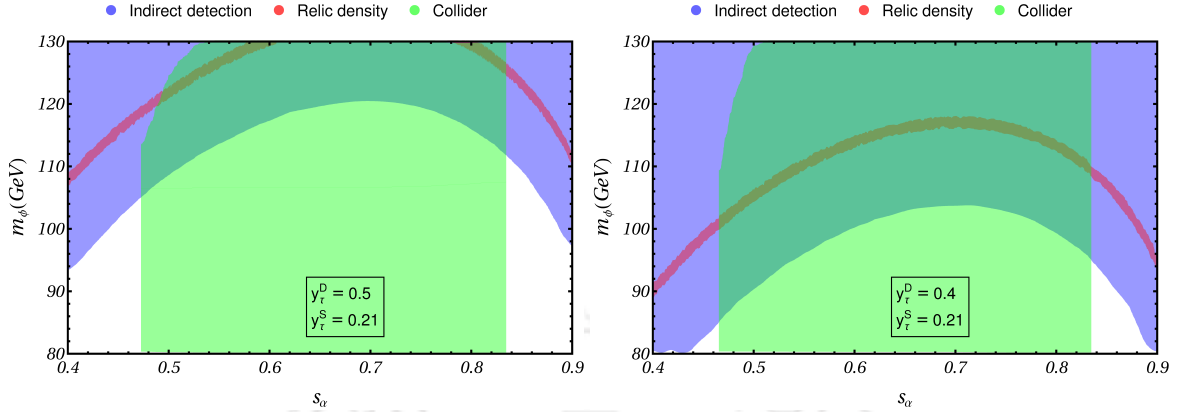


Figure 5.4.7: Parameter region combining all the experimental constraints. The blue shaded region denotes the region allowed by Fermi-LAT bounds, the red shaded area represents the relic density and the green shaded region gives the LHC limits for $3\tau + E_T^{miss}$ channel at $\mathcal{L} = 3 \text{ ab}^{-1}$, corresponding to $\mathcal{S} \geq 3$. Moreover, the entire parameter space is allowed by the direct detection constraints. The charged dark lepton masses are fixed at $(m_{\psi^\pm}, m_\chi) = (420, 150) \text{ GeV}$, whereas m_{ψ^0} varies with s_α and can be calculated from Eq. (5.2.10).

One can set limits on the DM mass and the mixing angle by observing the region of overlap between all the shaded areas. However, with proper choice of the couplings, the overlap can be tuned. We show this in the two plots in Fig. 5.4.7. Therefore, it can be concluded from these figures that the constraints on the mixing angle from the experimental limits can be tuned primarily with the proper choice of other parameters.

5.5 Summary

We have proposed a viable model with a leptophilic singlet scalar WIMP DM. A $\mathbb{Z}_2$ symmetry renders the stability to the dark sector. Apart from the DM candidate, the dark sector consists of a $SU(2)_L$ doublet and a singlet fermion. The presence of only a $SU(2)_L$ doublet in the dark sector which interacts through gauge as well as Yukawa couplings to the SM, adds new annihilation channels in the relic density calculation. However, in that case, one cannot distinguish one component of the doublet from the other by any means because of the degeneracy in their masses. Even if the degeneracy is lifted by the introduction of a new scale of EWSB through the extension of the scalar sector, both the doublet components interact through the same couplings and channels. This poses a serious problem in their distinction from one another. Introduction of a dark singlet in this scenario yields interesting features in the phenomenology. This is due to the fact that now, based on the electric charge of the new singlet, one component of the doublet will mix with it while the other component remains independent of mixing. Not only this additional degree of freedom automatically lifts the mass degeneracy of the dark fermions without making it an ad hoc proposition, but through the dark sector mixing, we can distinguish their effects in an experimental setup. Hence, on one hand, our model revives the simplest model of scalar singlet DM from the clutches of the stringent bounds of DM search experiments. On the other hand, it opens new search possibilities in the controlled environment of non-collider as well as collider based experiments. We conclude our observation and inferences in the following:

(1) For a better understanding of the mixing effects, in the relic density calculation, we choose the parameter region where the coannihilation channels are dominant. It is well known in the literature that the presence of coannihilation channels can boost up the relic density without adding to DM direct searches. For a similar model in Ref. [91] with leptophilic scalar DM and a fermionic doublet partner, we explored the viable parameter region thoroughly. In the present work, we show that in comparison with the previous study, here it is possible to relax the parameter region by a few orders of magnitude

for appropriate tuning of the mixing parameter. This is due to the fact that as the dark fermions can non-trivially coannihilate now, effectively more channels are added to the total annihilation cross section. To compensate for this increase, we showed in the analysis that for larger mixing in the coannihilation regime, one needs to have a smaller coupling in order to be relic density allowed. This makes a larger parameter space viable over the full range of the mixing compared to the previous work.

(2) We show in the analysis that mixing can be a very useful tool in discriminating between the dark sector particles of different isospins. This is because of the fact that for the two extrema of the mixing angle, one of the dark partners is purely a singlet while the other one remains a pure $SU(2)_L$ doublet. For the intermediate values of the mixings however, it is obvious that the charged dark partners are mixed states. For very low or high mixing, coannihilation of these two dark fermions with DM will be substantially different from each other. This primarily because of the channels involving W boson, which are available for the $SU(2)_L$ doublet dark partner for very low mixing and the singlet dark fermion for very high mixing. In our analysis, we discuss with correlation plots how these Gauge couplings help to clearly demarcate the parameter region w.r.t the contribution of the singlet and doublet dark partners towards total DM annihilation. However, as the mixing increases, the relative contribution of these dark sector particles accordingly vary. In the mixed scenario, it is interesting to observe how the viable parameter region evolves when the other parameters are fixed and it is only up to the mixing parameter to dictate the contribution of various dark sector coannihilation channels.

(3) The mixing can directly affect various DM search prospects, e.g., indirect detection and collider searches. In indirect detection, the velocity-averaged annihilation cross section has a dependence on mixing due to having mixed states in the propagator. We show in our analysis that for low DM mass, it is possible to relax the existing bound on the upper limit of Yukawa coupling in the presence of mixing.

(4) From the observations of the kinematic distribution of various observables for the $3\tau + E_T^{miss}$ and $\ell\tau + E_T^{miss}$ channels, we conclude that one can clearly distinguish the effects of the mixing parameter that remains unaffected by the change in other free parameters of the model. That this gradual change in the peak and tail positions of the kinematic distributions with the variations of mixing is independent of other free parameters, was further established by changing the mixing parameter and keeping the rest of the independent parameters fixed. We can ensure the presence of the mixing parameter between the dark sector particles of the theory by looking at the peak and tail positions these distributions. That this feature of mixing is not limited to the Dirac fermions only can be concluded from the other studies in the literature [104, 108].

(5) We have studied the statistical significance for the above two final states at 13 TeV LHC. Between the two, the $3\tau + E_T^{miss}$ channel can probe the signal for a larger range of mixing parameter. While it is not possible to observe the signal at the present LHC luminosity, we have shown with the numerical estimates that the signal can be observed at higher luminosities of 300 fb^{-1} and 3 ab^{-1} , which are projected to be achievable at LHC within a few years.

(6) Lastly, we have done a scan of the parameter space in DM mass vs mixing angle plane, combining all the experimental constraints coming from DM relic density, direct detection and indirect detection as well as the collider constraints. For fixed values of the dark sector masses, it is indeed possible to exclude a portion of parameter space, but this can be tuned with the proper choice of DM-SM couplings. This is due to the fact that the effects of the couplings and dark sector masses are non-trivial on the various phenomenological aspects of this feature-rich model. Therefore, it is not straightforward to obtain such exclusion limits on mixing angle without taking into account the effect of the other parameters.

We conclude this chapter with the assertion that the mixing between a singlet and a doublet dark sector fields can turn the table in favour of a quintessential scalar singlet DM model. It evades the stringent experimental bounds from the DM detection experiments as well as presents new opportunities for its detection in the lab.





Part III

Non-thermal dark matter



As a third alternative to standard WIMP scenario, in this part we discuss a completely different origin of DM, known as the freeze-in mechanism [154]. As a brief recap of freeze-in scenario in 1, unlike the WIMP case, here DM has very small coupling with SM and its abundance slowly builds up from zero through the decay or scattering of heavier thermal bath particles and saturates when DM/decaying particles become Boltzmann suppressed. Its feeble coupling with the visible sector keeps it safe from the direct searches. A natural consequence of non-thermal DM is the slow decaying partner particles that can leave distinct metastable/stable signatures [49] at the colliders due to their long-lived nature.

Here we have discussed a minimal model which can explain DM as both FIMP and WIMP depending on its coupling strength with thermal bath, i.e., SM as well as indicates possible long-lived signatures naturally arising from the requirement of the theory. Apart from providing novel charged LLP signatures, the model contains a long-lived right-handed neutrino as well, which originates naturally in the model arrangement. We propose methods of testing all these metastable particles in the present/future detectors.



Freeze-in and Freeze-out of Dark Matter with Charged Long-lived Partners

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6.1 Motivation

In this chapter, we have discussed non-thermal dark matter phenomenology as valid non-WIMP viability. The underlying mechanism, popularly known as *freeze-in mechanism* [154] as opposed to freeze-out of the WIMP, needs much smaller couplings, thus the particles are called Feebly Interacting Massive Particles (FIMP) [154]. In this case, the direct detection limits are not applicable, as the nuclear scattering cross sections are too small to be probed in the present or near-future experiments. However, it may be noticed that the FIMP mechanism works usually with a thermally produced partner particle decaying very slowly to the dark matter particle [69, 70, 126, 202, 218, 221]. Such feeble decays make these partner particles live longer than a typical unstable particle. For example, if produced in proton-proton collisions at the LHC, such long-lived particles (LLP) could propagate to a perceivable distance before they decay. It provides a possibility for a new way of identifying these scenarios at the collider experiments. Signatures typically include displaced vertex, disappearing charged tracks, and even kinks on the charged tracks. LHC is on the lookout for long-lived particles [2–8, 10, 11, 22, 178, 187, 210–214]. Apart from the dark matter scenarios discussed above, LLP appear in extra-dimensional models [36, 132], special scenarios in supersymmetry with very narrow mass splitting between the decaying particle and the final state particles [43, 44, 143, 156, 163, 200], SM extension with light neutrino mass [110, 125] and models explaining baryogenesis [100, 113]. The high-luminosity version of the LHC (HL-LHC) has dedicated upgrades keeping in view the possible signatures of the LLP [14, 103, 125, 139].

Besides, within the LHC complex, a dedicated detector is known as (MAssive Timing Hodoscope for Ultra Stable neutral pArticles) MATHUSLA [114, 191] is proposed as a large volume surface detector specifically looking for LLP. A detailed study of LLP signatures at the LHC in simplified FIMP models is carried out by [49, 120, 141], which shows that mass of the dark matter particles should be in the keV - MeV range to be compatible with dark matter relic density, at the same time providing decay lengths of the partner particles that could be probed by LHC.

We consider a novel scenario where a fermionic dark matter (ψ) is produced from the slow decay of a charged fermion (χ^+). This necessitates the presence of a charged bosonic degree of freedom, which we consider as a charged scalar S^+ for simplicity, allowing χ^+ to decay through $\chi^+ \rightarrow \psi S^+$. The Yukawa coupling corresponding to this interaction can be adjusted to obtain the required relic abundance of the dark matter. Considered as an $SU(2)_L$ singlet, the charged scalar requires the presence of a gauge singlet neutral fermion N for it to decay via $S^+ \rightarrow N\ell^+$ channel. The Yukawa coupling corresponding to this interaction is independent of the dark sector dynamics. In particular, it is possible that it is small enough to make the charged scalar have sufficiently large decay length to leave distinct LLP signature at the LHC. On the other hand, if the Yukawa coupling in the dark sector is large enough to accommodate thermal production of the dark matter, we have the WIMP scenario with freeze-out relic density. However, one should be careful in dealing with large Yukawa couplings as, in the absence of tree-level diagrams, they can still contribute to the direct detection cross section through photon/ Z boson mediated penguin diagrams.

Typically, one may imagine that when thermal production is enabled, the non-thermal production through the charged fermion decay is not significant. However, an interesting scenario arises when the non-thermal production adds significantly to the number density along with the thermal production. Such a scenario would emerge when the relevant couplings and masses are in favourable ranges. The relic density, in this case, will be dictated by the thermal freeze-out mechanism. Nevertheless, this possibility will lead to more viable regions of parameter space, with consequent effects in other sectors including the collider processes. The result of all these effects provides a feasible mechanism of dark matter of mass in a wide range of a few GeV to close to a TeV, along with LLP that could be probed through the LHC experiments. Presence of multiple LLPs potentially provides another interesting aspect to their search at the LHC through cascade decay of $LLP \rightarrow LLP \rightarrow SM$, adding a different dimension to the LLP searches. In this article, we shall study the details of the dark matter scenario of the model and indicate the collider possibilities, with the detailed collider phenomenology deferred for a future study.

6.2 Model details

We consider the dark matter to be a gauge singlet fermion denoted by ψ , which is made stable by imposing a $\mathbb{Z}_2$ symmetry under which it is odd. The non-thermal production of ψ is realised through the slow decay of a partner charged fermion χ^+ , which is a vector-like gauge singlet. χ^+ and ψ have the same transformation property under the $\mathbb{Z}_2$. This decay necessitates the presence of a charged gauge singlet scalar field, denoted here by S^+ , which in turn decays to the charged standard lepton and newly introduced gauge singlet vector-like neutral fermion, N . The Yukawa interaction of N with the Standard

	S^+	N_1, N_2, N_3	χ^+	ψ
Y	+2	0	+2	0
$\mathbb{Z}_2$	+	+	−	−

Table 6.2.1: Additional fields and their charges. Electromagnetic charges are given by $Q = t^3 + Y/2$.

Model (SM) Higgs field and the left-handed lepton doublet along with its explicit Dirac mass term, could be made use of to generate the neutrino mass through the type-I seesaw mechanism. Keeping this in

mind, we introduce three copies of N in our set up. The additional particle spectrum along with their hypercharge and $\mathbb{Z}_2$ charge are given in Table 6.2.1. All the SM fields are even under the $\mathbb{Z}_2$.

With the above particle content, the Lagrangian of the model is given by

$$\begin{aligned}\mathcal{L} = & \mathcal{L}_{\text{SM}} + (D_\mu S)^\dagger (D^\mu S) + \bar{\chi} i \gamma^\mu D_\mu \chi + \bar{\psi} i \gamma^\mu \partial_\mu \psi + \sum_i \bar{N}_i i \gamma^\mu \partial_\mu N_i \\ & - m_\chi \bar{\chi} \chi - m_\psi \bar{\psi} \psi - \sum_i m_{N_i} \bar{N}_i N_i \\ & - \left(y_1 \bar{\chi} S \psi + \sum_i y_{2i} \bar{N}_i S \ell_i + \sum_i y_{N_i} \bar{L}_i \tilde{\Phi} N_i + \text{h.c.} \right) \\ & - (\mu_S^2 S^\dagger S + \lambda (S^\dagger S)^2 + \lambda_1 S^\dagger S \Phi^\dagger \Phi)\end{aligned}\quad (6.2.1)$$

where Φ represents the SM Higgs doublet, L_i and ℓ_i denote the SM left-handed lepton doublet and right-handed charged lepton singlet, respectively. The summation index i runs from 1 to 3 indicating the three flavours of leptons. The Yukawa couplings and the mass parameter m_{N_i} are taken to be diagonal so that the interactions do not lead to any flavour-violating processes. However, the contribution to $g - 2$ of the charged leptons arising at the one-loop level is proportional to y_{2i}^2 , and therefore require to be small [168, 225] for the first two generations, while the $i = 3$ it could be more relaxed. Moreover, the presence of neutral fermions N_i can explain neutrino mass generated via Type-I seesaw mechanism. The mass matrix in the interaction basis is

$$\begin{pmatrix} 0 & v Y_N \\ v Y_N^T & m_N \end{pmatrix}, \quad (6.2.2)$$

where Y_N and m_N are 3×3 matrices with $Y_N = \text{diag}(y_{N_1}, y_{N_2}, y_{N_3})$ and $m_N = \text{diag}(m_{N_1}, m_{N_2}, m_{N_3})$. Keeping m_{N_i} at the GeV-TeV scale requires $y_{N_i} \sim 10^{-8}$ to get the right neutrino mass (of < 0.1 eV). The other Yukawa coupling is limited by perturbativity, and thus we consider $y_1 \leq \sqrt{4\pi}$. Electroweak symmetry breaking is controlled purely by the Higgs mechanism involving the Φ field. After the symmetry breaking, the mass of the charged scalar is given by

$$m_S^2 = \mu_S^2 + \frac{\lambda_1 v^2}{2}, \quad (6.2.3)$$

whereas m_χ and m_ψ are free parameters. In addition to the constraints discussed above, further constraints come from the collider experiments. If kinematically allowed, $\chi^+ \chi^-$ and $S^+ S^-$ pairs could be produced in $e^+ e^-$ collision, leading to dilepton final state accompanied by missing energy. Consequently, to avoid LEP constraints we consider $m_\chi \geq 45$ GeV. Similarly, the Higgs bosons can decay to $S^+ S^-$ pair owing to the quartic coupling, λ_1 , contributing to the Higgs decay. For simplicity, we consider $m_S \geq 63.5$ GeV to kinematically prohibit this decay. On the other hand, possible contributions to $H \rightarrow \gamma\gamma$ and to $H \rightarrow \psi\psi$ through the presence of S^+ and χ^+ in the loop may need attention. We discuss the contribution to the di-photon decay channel in Section 6.B. With $m_S = 150$ GeV, this contribution (Eq. (6.B.7)) leads to less than 3% contribution for $\lambda_1 = 0.01$. Contribution to the invisible Higgs decay is negligible compared to the standard contributions, as discussed in Section 6.C.

Coming to other constraints, the decay pattern of $S \rightarrow N \ell$ could look very similar to supersymmetric slepton decays to neutralinos and leptons, and therefore it is tempting to imagine that the collider constraints from slepton search apply to S^+ . However, the SUSY searches are quite model dependent, and thus not applicable as such in our case. For example, constraints on long-lived stau quoted in Ref. [107] is derived within the restricted mass-splitting scenario and within the context of gauge mediated SUSY breaking. Since we have the freedom to choose the couplings and masses within a much larger range, to start with, we do not consider any such collider constraints. However, in the full considerations of collider phenomenology, these constraints should be included.

Turning to the dark matter side, we have different possibilities arising here depending on the kinematics and couplings. First of all, the dominant mechanism for the production of the dark matter ψ is

the decay of thermally produced charged fermionic partner, χ^+ . This immediately requires $m_\chi > m_\psi$. The decay proceeds through their coupling with the singlet charged scalar, S^+ . Here, there are two possibilities: (i) when kinematically allowed with $m_\chi \geq m_\psi + m_S$, the decay essentially goes through the on-shell S^+ ($\chi^+ \rightarrow \psi S^+$); (ii) when $m_\psi + m_N \leq m_\chi < m_\psi + m_S$, it results in a three-body decay of $\chi^+ \rightarrow \psi S^* \rightarrow \psi N \ell^+$. In the first case, the decay is dictated by a single coupling y_1 , whereas in the second case it is the product of $y_1 y_2$ that matters. As noted above, y_2 is restricted to be small for the first two generations, while it can be more relaxed for the third generation which couples the τ lepton with S^+ . Direct detection experiments can constrain the coupling y_1 through processes at the one-loop level. We have discussed the contribution in the Section 6.D. To have $\sigma_{\text{SI}}^{\psi\chi} \leq 10^{-46} \text{ cm}^2$ as required by the direct detection experiments, for masses of the order of 100 GeV, the coupling typically comes around $y_1 \lesssim 0.05$.

The strength of all these interactions will dictate the relic density of ψ , the dark matter candidate. In Table 6.2.2 we have listed the distinct possibilities that could arise.

Case 1 : $m_\chi > m_S + m_\psi$				
Mechanism	y_1	m_ψ	Possibilities	Viable?
Thermal production	large	$< m_S$	over-abundance	No
		$> m_S$	$\psi\psi \rightarrow S^+ S^-$ enabled freeze-out	Yes
Production through decay of χ^+	small	—	freeze-in to saturation	Yes

Case 2 : $m_\psi + m_N < m_\chi < m_S + m_\psi$					
Mechanism	$y_1 \ y_2$	y_1	m_ψ	Possibilities	Viable?
Thermal production	large	any	$< m_S$	over-abundance	No
		restricted	$> m_S$	$\psi\psi \rightarrow S^+ S^-$ is enabled	Yes
Thermal production + through decay of χ^+	small	small	any	freeze-in to saturation	Yes
		restricted	$> m_S$	$\psi\psi \rightarrow S^+ S^-$ enabled freeze-out	Yes
			$< m_S$	over-abundance	No

Table 6.2.2: Distinct scenarios and their viabilities considering relic density of the dark matter.

In the first case with 2-body decay of χ^+ , the coupling y_2 does not affect the scenario. For large values of y_1 , the dark matter ψ is thermally produced. In addition to the decay of χ^+ , the scattering process $S^+ S^- \leftrightarrow \chi^+ \chi^-$ will also be relevant here. The only way to get the right relic density is by taming the dark matter density through its annihilation. This could be achieved in the present model through the pair annihilation of ψ into $S^+ S^-$ pairs and possible coannihilations when kinematically allowed. Note that it is the same coupling which dictates the annihilation process as well. On the other hand, for very small y_1 , the slow production of ψ gradually builds up the dark matter density. Apart from y_1 , the mass splittings would also play a crucial role in getting the right relic density. Moving to the case of 3-body decay of χ^+ as listed in Table 6.2.2 Case 2, it is the product of the couplings $y_1 y_2$ which is now relevant. Again, for large values of this coupling combination, ψ is thermally produced, and the over-abundance could be tamed with the pair annihilation of ψ for favourable kinematic ($m_\psi > m_S$) and dynamic ($y_1 > 0.1$, as shown through numerical studies discussed later) conditions. This is quite similar to the large coupling scenario of case 1. On the other hand, for very small values of $y_1 y_2$, there

are two different situations. One is that y_1 is also very small. In that case, the only relevant mechanism is the freeze-in mechanism. On the other hand, with large enough y_1 thermal production can lead to overabundance when $m_\psi < m_S$, which is tamed through the annihilation process, $\psi\psi \rightarrow SS$ for $m_\psi > m_S$. In this scenario, when $y_1 y_2$ is not too small, significant contribution from non-thermal (decay of χ) production of ψ can add to the relic density. However, proper treatment of this requires solving the coupled Boltzmann equation for both ψ and S , as ψ and S may no more be in chemical equilibrium. In the present discussion, we shall restrict ourselves to the two limiting cases of either the thermal production or non-thermal production dominate, with the other contribution negligible.

In the following sections, we shall discuss the freeze-in and freeze-out scenarios separately, and present a numerical study to understand the specific situations and viable parameter space regions in each of these cases.

6.3 Freeze-in solution

The production of DM candidate ψ through the slow decay of gauge produced χ^+ gives rise to the freeze-in mechanism by building the DM density to the required value. The decay of χ^+ is controlled by the Yukawa coupling y_1 with S^+ and ψ . When kinematically possible, it will be a two-body decay with S^+ produced on-shell. On the other hand, if $m_S + m_\psi \geq m_\chi$, it is a three-body decay to final state $\psi N \ell$, mediated by a virtual S^+ . In the latter case, the decay is dictated by the Yukawa coupling y_2 of $S^+ N \ell^+$ along with y_1 . While in the first case y_1 is required to be small enough for freeze-in to be possible, in the second case the product $y_2 y_1$ should be of this order. The relic density depends on the decay width, Γ_{χ^+} and the masses of the parent particle and the DM candidate through the relation [154]

$$\Omega_\psi h^2 = \frac{2.19 \times 10^{27} g_\chi}{g_*^S \sqrt{g_*^\rho}} \frac{m_\psi \Gamma_{\chi^+}}{m_\chi^2}, \quad (6.3.1)$$

where $g_*^{S,\rho}$ are effective degrees of freedom in the bath at freeze-in temperature $T \sim m_\chi$. g_χ denotes the number of degrees of freedom corresponding to the decaying particle χ^+ , which is equal to 2 for a fermion.

We shall consider these two distinct cases separately in our numerical analysis below. The charged scalar decays to $N \ell^+$, enabled by the Yukawa coupling y_{2i} . In the first case of two-body decay of χ^+ , the decay of S^+ is decoupled from the freeze-in mechanism, and therefore y_2 can be independently varied. As mentioned in the introduction, this coupling can potentially induce anomalous $g - 2$ for the charged SM leptons, and therefore y_2 is required to be small. With small y_2 , S^+ may have delayed decay, but within the detectors of the LHC experiments. The heavy neutrino N can decay through its coupling with the SM neutrino and the Higgs boson, y_N , which is required to be very small to generate the right neutrino mass. Thus, potentially, N could also be long-lived and mostly decay outside the detector complex giving missing energy signature. In the second case of the three-body decay of χ^+ , although not decoupled from the relic density formation, y_2 can be such that S^+ has delayed decay to leave distinct signature at the collider detectors. We shall discuss the signatures of this long-lived S^+ in the later section. Below we shall focus on the dark matter side of the model and find viable parameter space regions compatible with observations. Feynman diagrams corresponding to the decay of χ^+ in the two cases, and that of the S^+ decay are given in Figs. 6.3.1a)–c), respectively. We shall discuss these two cases separately below.

6.3.1 Two-body decay of χ^+ with $m_\chi > m_S + m_\psi$

In this case, the DM candidate is produced through the two-body decay of the χ^+ which is gauge produced abundantly, and in thermal equilibrium. The decay width of χ^+ is given in this case by,

$$\Gamma_{\chi^+ \rightarrow S^+ \psi} = \frac{y_1^2}{16\pi m_\chi^3} [(m_\chi + m_\psi)^2 - m_S^2] [(m_\chi^2 - m_S^2 - m_\psi^2)^2 - 4m_S^2 m_\psi^2]^{1/2} \quad (6.3.2)$$

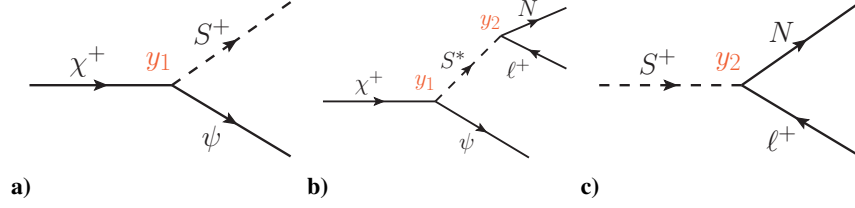


Figure 6.3.1: Feynman diagram showing the two body (a) and three body (b) decays of χ^+ , and the Feynman diagram showing the decay of S^+ (c).

One may notice that, in principle, a pair of ψ can annihilate into $S^+ S^-$ pairs when kinematically feasible. Such a reaction will be a t -channel process mediated by χ^+ as shown in the Feynman diagram in Fig. 6.3.2(a). In addition to the pair annihilation, when the mass splitting is favourable, the co-annihilation processes with Feynman diagrams shown in Fig. 6.3.2(b)-(d) could also affect the number density of the dark matter. However, the cross section of the pair annihilation process is proportional to y_1^4 , while that of all other processes are proportional to y_1^2 , the same coupling that controls the slow production leading to the freeze-in mechanism. Thus, the annihilation cross sections are expected to be negligibly small leaving no signature in the relic density. The relic density for a wide range of parameters

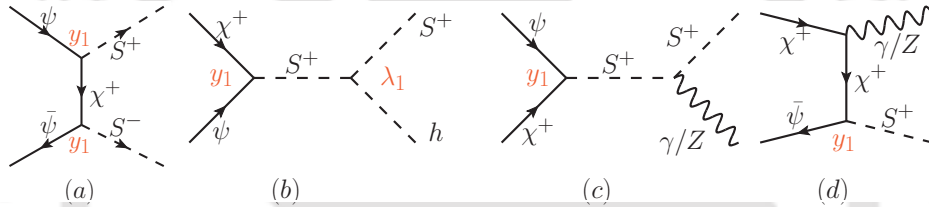


Figure 6.3.2: Pair annihilation (a) and co-annihilation channels (b, c, d) of the dark matter ψ .

is computed using MATHEMATICA with the help of the above relation and compared with the value allowed by PLANCK [16] experiment. The relic density, in this case, depends on the single Yukawa coupling y_1 , and the masses of the participating particles, m_χ , m_S and the dark matter particle m_ψ . We fix the other parameters, which do not affect the relic density as $\lambda = \lambda_1 = 0.01$, $y_N = 10^{-8}$, $y_2 = 10^{-9}$.

All the three copies of the neutral vector fermion mass are fixed at $m_N = 30$ GeV in our scan. While the decay width of S^+ has a kinematic dependence on m_N , the width of χ^+ itself is not affected by this parameter. The range of the relevant parameters used in the scan is given in Table 6.3.1. All through we have maintained the mass splitting so that χ^+ can undergo a two-body decay. Scanning over 50,00 random points, we selected those satisfying the right relic density and analysed those for possible correlations and constraints.

$1 \text{ GeV} \leq m_\psi \leq 1 \text{ TeV}$	$65 \text{ GeV} \leq m_S \leq 1 \text{ TeV}$
$m_S + m_\psi \leq m_\chi \leq 1 \text{ TeV}$	$10^{-14} \leq y_1 \leq 10^{-8}$

Table 6.3.1: Range of relevant parameters considered for the numerical scan in this study.

In Fig. 6.3.3 the correlation between y_1 and relevant masses are given. As seen from y_1 vs m_S larger masses require larger couplings to compensate for the effect of the phase space. On the other hand, in the case of y_1 vs m_ψ , the phase space suppression for larger m_ψ is smaller compared to the enhancement in the invariant amplitude, as is clear from the decay width expression in Eq. (6.3.2). Further, there is an additional factor of m_ψ in the expression for $\Omega_\psi h^2$ in Eq. (6.3.1), pulling down the decay width to maintain the same relic density. The correlation between the parent particle mass (m_χ) and y_1 indicates that y_1 is preferred to be in the range of $10^{-13} - 10^{-9}$ to satisfy the required relic density. As expected, smaller mass splitting $\Delta m_{\chi S} = m_\chi - m_S$ requires larger y_1 to compensate for the narrow phase space

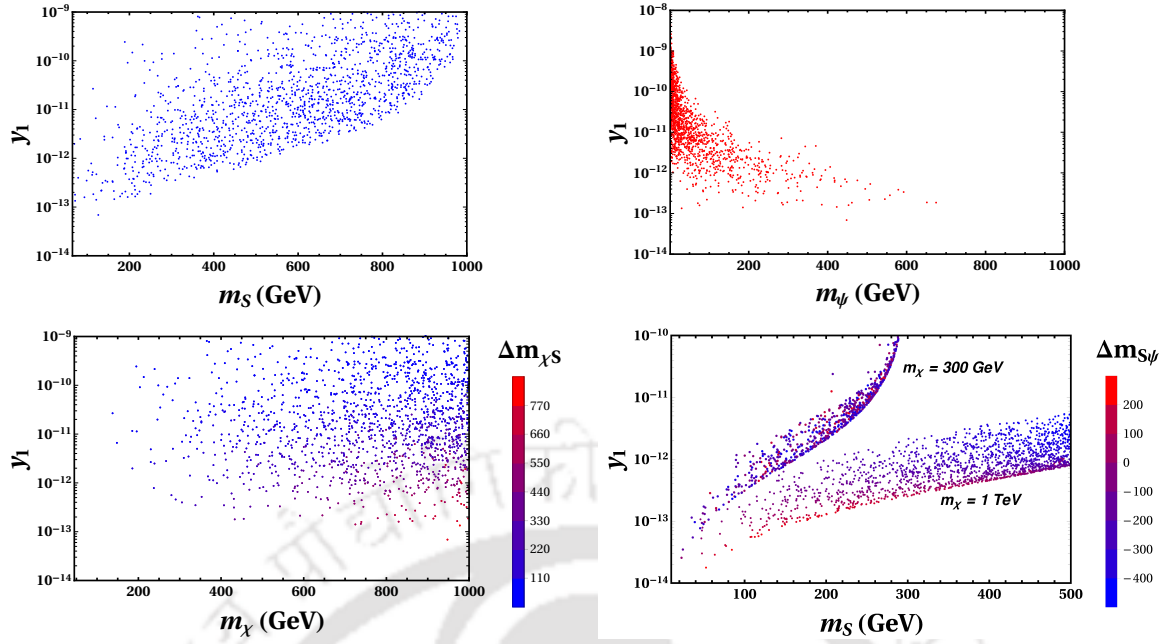


Figure 6.3.3: Yukawa coupling y_1 vs mass of the non-SM particles that satisfy the relic density bound. In plots in the top row and bottom-left m_S , m_ψ and m_χ are varied from a few GeV to 1 TeV. Other parameters are as per Table 6.3.1. In the bottom-right plot, $m_\chi = 300$ GeV and 1 TeV are considered, while m_ψ is varied in the full range.

available. The right-bottom plot in Fig. 6.3.3 shows y_1 against m_S for two different m_χ values of 300 GeV and 1 TeV. The mass splitting $\Delta m_{S\psi} = m_\psi - m_S$ is indicated in different colours. While in the case of lighter m_χ the mass splitting is not very relevant, in the case of heavier χ^+ smaller y_1 is preferred for larger m_ψ values, in agreement with top-right plot and the discussion above.

6.3.2 Three body decay of χ^+ with $m_\psi + m_N + m_\ell < m_\chi < m_S + m_\psi$

In this case, the decay of χ^+ to on-shell S^+ is not kinematically possible. However, three body decay enabled through a virtual S^+ decaying into $N\ell$ pair, as illustrated in the Feynman diagram in Fig. 6.3.1 (b), produces the dark matter. The decay width $\Gamma_{\chi^+ \rightarrow \psi N \ell^+}$ is proportional to $(y_1 y_2)^2$, as given by

$$\Gamma_{\chi^+ \rightarrow \psi N \ell^+} = \frac{(y_1 y_2)^2}{256\pi^3 m_\chi^3} \int_{X_{\min}}^{X_{\max}} dX [(m_\chi^2 - X - m_\psi^2)^2 - 4X m_\psi^2]^{1/2} [(X - m_N^2 + m_\ell^2)^2 - 4X m_\ell^2]^{1/2} \times \frac{[(m_\chi + m_\psi)^2 - X][X - m_N^2 - m_\ell^2]}{X[(X - m_S^2)^2 + \Gamma_S^2 m_S^2]}, \quad (6.3.3)$$

where $X_{\max} = (m_\chi - m_\psi)^2$ and $X_{\min} = (m_N + m_\ell)^2$. Having the product of the couplings controlling the production, unlike in the previous case, the following distinct possibilities arise in this case, depending on the mass splitting $\Delta m_{\psi S} = m_\psi - m_S$.

1. $y_1 y_2 \ll 1$, $y_1 \ll 1$ for any $\Delta m_{\psi S}$: In this case, the coupling y_1 being small ($\leq 10^{-8}$), the annihilation channel $\psi\psi \rightarrow S^+ S^-$ is not effective leading to a situation similar to Case 1 described in Section 6.3.1. ψ produced via 3-body decay of $\chi^+ \rightarrow \psi N \ell^+$ slowly saturates to leave the required relic density. We shall show that, for a large range of mass of S^+ , it is possible to find couplings compatible with the freeze-in scenario in the range of $10^{-11} \leq (y_1 y_2) \leq 10^{-5}$.
2. $y_1 y_2 \ll 1$, $y_1 \sim 1$, with $\Delta m_{\psi S} < 0$: In this case the production is mostly thermal through $SS \rightarrow \psi\psi$ with possible addition from the decay of χ^+ , the contribution being dependent of the value of $y_1 y_2$. Since the annihilation $\psi\psi \rightarrow S^+ S^-$ is kinematically disfavoured, it results in

over-abundance. The co-annihilation channels could, however, be relevant, as the kinematics may not restrict those.

3. $y_1 y_2 \ll 1$, $10^{-8} < y_1 \sim 1$, with $\Delta m_{\psi S} > 0$: In this case again, the thermal production through $SS \rightarrow \psi\psi$ and from the decay of χ^+ builds up the number density of the dark matter particle. However, owing to sizeable y_1 the now kinematically allowed annihilation cross section is significant, leading to favourable relic density through freeze-out. The resultant relic density at the saturation depends on various factors including the rate of decay of χ^+ and the annihilation cross section of $\psi\psi \rightarrow S^+ S^-$ channel.

We shall not discuss case 2 above any further, as this does not lead to the right relic density. The other two cases, case 1 where the thermal production and annihilation are not relevant, owing to the small value of y_1 , and case 3 with large y_1 are discussed with numerical results below.

In the absence of annihilation : We shall first consider the situations 1 listed above, where the annihilation does not play a role. Scanning the parameters within the range specified by Table 6.3.2, we explore the parameter space providing the right relic density.

$1 \text{ GeV} \leq m_\psi \leq 1 \text{ TeV}$	$65 \text{ GeV} \leq m_S \leq 1 \text{ TeV}$
$m_N + m_\psi \leq m_\chi \leq 1 \text{ TeV}$	$10^{-12} \leq y_1 y_2 \leq 10^{-5}$

Table 6.3.2: Range of relevant parameters considered for scan. The condition, $m_\chi < m_\psi + m_S$ to kinematically disable $\chi^+ \rightarrow S^+ \psi$ with S^+ on-shell is further imposed. A few selected m_N values are considered for the analysis.

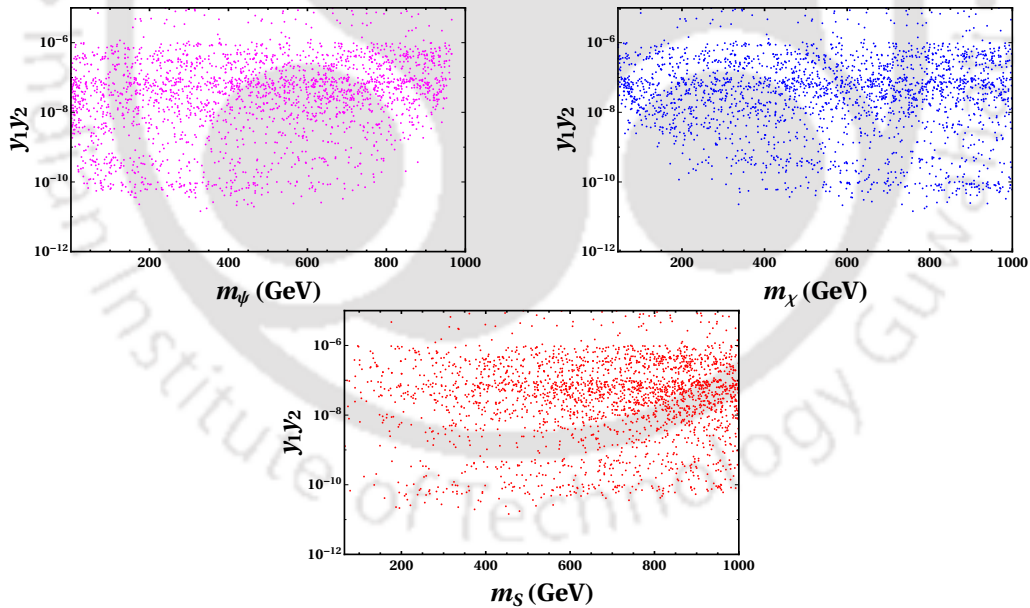


Figure 6.3.4: Correlation between the relevant coupling and the dark matter mass m_ψ (left), the decaying parent particle mass m_χ (right) and the associated charged scalar S^+ (bottom), satisfying the relic density bound.

In Fig. 6.3.4 the coupling combination $y_1 y_2$ is plotted against m_ψ (top left figure), m_χ (top right figure) and m_S (bottom figure). Generically, the product of the coupling below 10^{-10} is not possible, with a few points between 10^{-11} and 10^{-10} corresponding to relatively heavier χ^+ and lighter m_S so that the propagator factor compensates for the small coupling. For lighter χ^+ below 200 GeV, the coupling is mostly larger than 5×10^{-10} , which is filtered down to slightly lower than 10^{-10} for masses between 200 and 400 GeV, above which it is possible to go down by one more order. The correlation between m_χ and m_ψ is plotted in Fig. 6.3.5, where the effect of m_N is indicated with two illustrative

values of $m_N = 30$ and 200 GeV. The effect mostly arises from the kinematic limit constraint of $m_\chi > m_\psi + m_N$. In the right-hand-side plot in Fig. 6.3.5, the effect of the coupling on this correlation is indicated for three different choices of couplings, for a fixed $m_N = 30$ GeV. Note that the decay width depends on the inverse powers of its mass, and therefore, require larger values of the couplings for heavier χ^+ , as clearly indicated in the three cases presented in the plot.

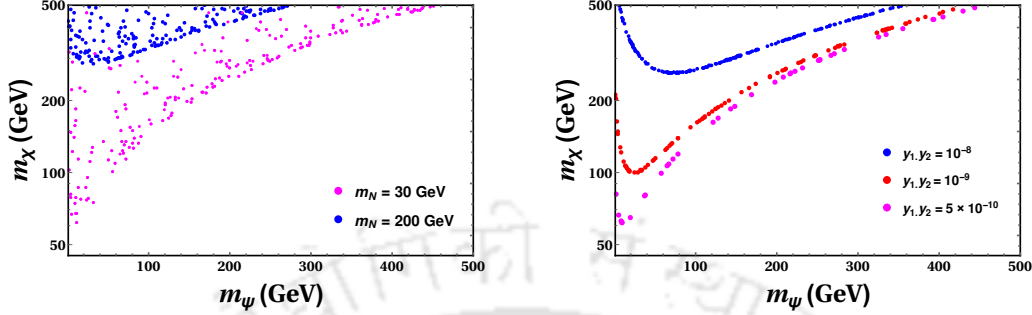


Figure 6.3.5: Correlation between m_ψ and m_χ in the case of freeze-in with 3-body decay of χ^+ . The *left* plot allows m_S and $y_1 y_2$ to vary within the specified range, while the *right* plot has $m_S = 800$ GeV and three different indicative $y_1 y_2$ values. m_N is fixed at 30 GeV in the *right* plot.

In the presence of annihilation : The third situation listed above, which includes the effect of sizable annihilation cross sections is interesting, however not usually considered in the literature. In most situations studied, a single coupling parametrised the interactions. Freeze-in mechanism requiring this to be small naturally makes the effects of annihilation irrelevant. However, as we demonstrate here, in the presence of an additional coupling such as present in the scenario discussed in this article provides a natural framework where along with the thermal production, a significant amount of DM is produced through the decay of the associated particle. In presence of annihilation, this leads to an interesting interplay of the thermal and non-thermal production of ψ . When y_1 and $y_1 y_2$ both are small, ψ is not in equilibrium with the thermal bath and is only produced from the decay of χ . However, S^+ , being gauge produced, if y_1 increases, $\psi\psi \rightarrow SS$ forward and backward scatterings also produce ψ from the thermal bath and slowly drive it towards equilibrium in the early Universe. This feature is illustrated in Fig. 6.3.6, where the yield of ψ is plotted against m_ψ for a fixed $y_1 y_2$ but for different y_1 values. The green curve can be considered as the typical non-thermal production via freeze-in due to the minuscule $\langle\sigma v\rangle$ (corresponding to $y_1 \sim 10^{-8}$). The blue and the red line correspond to larger y_1 which clearly show the increasing effect of the thermal production of ψ from the bath in the early Universe. However, for such values of $\langle\sigma v\rangle$, the non-thermal decay is still the dominant production channel and we see the effect of saturation of yield when it becomes Boltzmann suppressed, similar to the freeze-in scenario. Now, on increasing y_1 further, the thermal production becomes more dominant and the typical freeze-out scenario is recovered, which can be seen from the brown, pink, magenta and the black line. Similar to the WIMP freeze-out, these cases give overabundance or underabundance depending on the $\langle\sigma v\rangle$ value considered.

A few benchmark points compatible with this scenario are presented in Table 6.3.3¹, indicating the cases when pair annihilation and co-annihilation channels contribute significantly. For clarity, we have included the respective cases with small y_1 , where the effect of annihilation is absent, and the correct relic density is not achievable. While the results are presented for a fixed $m_N = 30$ GeV, we have checked that it produces the relic density within the same order of magnitude for a range of m_N from 1 GeV to a couple of 100 GeV. It is clear from Fig. 6.3.6 that based on the couplings, the DM is either

¹In Fig. 6.3.6 and Table 6.3.3, large values of y_1 is taken to show the effect of the interplay between the thermal and non-thermal production of ψ . However, direct detection diagrams in one loop constrains this coupling to $y_1 \lesssim 0.05$. This implies that this coupling in the allowed limit is not sufficient to satisfy the observed relic density, but it can make the thermal production dominant over the non-thermal case and lead to overabundance. We will show later on in the discussion of freezeout that coannihilation and mediator annihilation channels can satisfy the observed relic even with small y_1 .

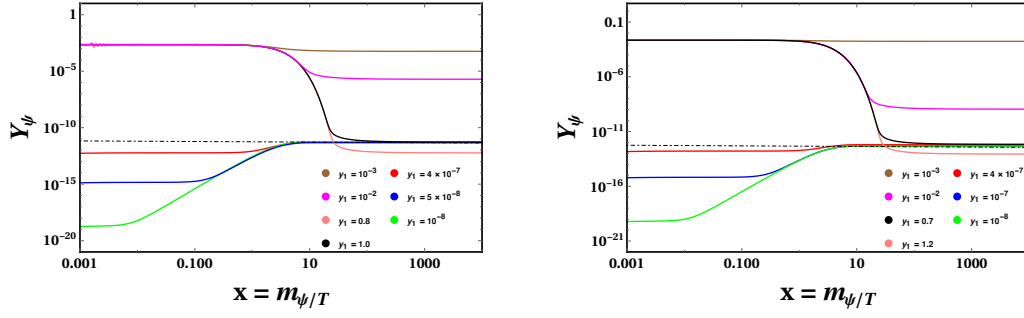


Figure 6.3.6: Plots showing interplay of the thermal and non-thermal production of ψ for various values of y_1 . The non-thermal production of ψ from χ decay is kept fixed with $y_1 y_2 = 10^{-10}$ (left) and $y_1 y_2 = 10^{-6}$ (right). The dot-dashed lines represent the correct relic density for the specific benchmark point $m_S=50$ GeV, $m_\psi=80$ GeV, $m_\chi=120$ GeV in the left plot and $m_S=463$ GeV, $m_\psi=941$ GeV, $m_\chi=972$ GeV in the right plot. m_N is fixed at 30 GeV.

m_χ	m_ψ	m_S	$y_1 y_2$	y_1	Ωh^2
515.5	485.4	176.3	4.7×10^{-6}	10^{-8}	0.119
				10^{-2}	2.1×10^5
				0.5	0.12
				1.0	0.008
394.5	352.9	338.5	1.1×10^{-8}	10^{-8}	0.119
				10^{-2}	2.2×10^6
				0.98	0.12
				1.5	0.02
460.6	300.2	214.1	3.72×10^{-11}	10^{-8}	0.119
				10^{-2}	5.9×10^5
				0.67	0.12
				1.0	0.02

Table 6.3.3: Benchmark points showing the effect of the additional thermal production of ψ with normal non-thermal production from χ^+ decay. For negligible y_1 , the correct relic density is achieved by non-thermal production only. Increasing y_1 enhances the thermal production of ψ through $SS \leftrightarrow \psi\psi$ channel in the early Universe, suppressing the non-thermal production. Therefore, the freeze-out scenario is recovered, where overabundance or underabundance is dictated by the choice of y_1 . All masses are in GeV. The mass of the neutral fermion is set to $m_N = 30$ GeV in all cases.

following freeze-in saturation or is driven towards thermal freezeout. This is largely due to the coupling y_1 which, for large values, is responsible for sufficient energy exchange between S^+ and ψ in the early Universe, therefore, suppressing the non-thermal production of ψ . If the production and annihilation channels of ψ could be considered independently, then it would be possible to restrict its production to only non-thermal decay of χ and annihilate through some other coupling which does not contribute to the energy exchange of DM with a thermal bath. A possible way out is to consider cases like mediator-driven annihilation and conversion as the possible modes of the annihilation of ψ . It is well known that for very small mass splitting between ψ and χ , $\chi\chi \rightarrow \text{SM SM}$ processes contribute substantially to the relic density of ψ . In this model, these processes are gauge mediated, hence independent of y_1 . Hence, even if we take both y_1 and y_2 small enough to ensure only non-thermal decay of χ^+ , the production and the annihilation of ψ remain independent of each other. Ref. [172] discusses a similar situation and on solving out-of-equilibrium Boltzmann equations (see Eqn. 9 and 10 of Ref. [172])

for DM candidate and the mediator, an interesting situation is obtained which can be interpreted as an intermediate stage between freeze-in and freeze-out scenario, where the non-thermal overproduction can be tamed by sufficient annihilation. However, as seen from the parameter choice of this plot, this is valid only for a small mass splitting between the DM candidate and the mediator, which is not of much importance for our discussion since the viable parameter space is very much contrived.

6.4 Exploring freeze-out possibility

The model can also explain thermal dark matter, such as the standard WIMP scenario through thermal freeze-out, which is viable over a large region of parameter space. In this case, it is assumed that the DM candidate ψ is already in equilibrium with the thermal bath. Hence we no longer require the DM interaction coupling y_1 to be very small unlike freeze-in scenario, therefore it is considered to be of the order of unity. All the possible annihilation diagrams are given in Fig. 6.3.2. Based on the mass splitting between the dark sector particles (χ^+ and ψ), the relic density will be dominated by either annihilation of DM into SM particles (Region 1) or co-annihilation between the dark fermions (Region 2). Fig. 6.3.2(a) represents the only diagram possible when ψ pair annihilates into $S^+ S^-$ pairs through t -channel mediation of χ^+ . Fig. 6.3.2(b)-(d) represent the co-annihilation channels which become effective only when the mass splitting between incoming particles is very small. The model is implemented in FEYNRULES [23] and the freeze-out relic density computation is done using MICROMEGAS [50].

The parameter space is classified into two regions of $m_\chi \gg m_\psi$ and $m_\chi \sim m_\psi$, with the co-annihilation channels becoming important in the latter case. We shall discuss these two regions separately in the following subsections.

6.4.1 Region with large mass splitting: $m_\chi \gg m_\psi$

The only pair annihilation of the dark matter in this region proceeds through the t -channel to a final state of $S^+ S^-$ pairs shown in Fig. 6.3.2(a), with the cross section proportional to y_1^4 . The co-annihilation process has a negligible contribution as the cross section for such process goes like the exponential of mass splitting between the dark matter and the partner, $e^{-(m_\chi - m_\psi)/m_\psi}$ [148]. To study the relic compatible parameter space regions, we scan over the available/relevant range of parameters and compute the relic density using micrOMEGAS. As in the earlier case, the irrelevant quartic couplings, λ and λ_1 are fixed at 0.01, and the Yukawa coupling dictating the neutrino mass are set at $y_{N_i} \approx 10^{-8}$. The mass of the neutral fermions is set to $m_N = 30$ GeV, while the Yukawa coupling between $S^+ N \ell$ is taken as $y_2 = 10^{-6}$. The parameters, which are relevant to the annihilation process here are varied randomly in the range given in Table 6.4.1.

$1 \text{ GeV} \leq m_\psi \leq 1 \text{ TeV}$	$m_\chi = m_\psi + (100, 800) \text{ GeV}$
$65 \text{ GeV} \leq m_S \leq 1 \text{ TeV}$	$0.0 \leq y_1 \leq 3.0$

Table 6.4.1: Range of parameters considered for the scan in Region 1: $m_\chi \gg m_\psi$ with thermal freeze-out of dark matter ψ .

In Fig. 6.4.1, the variation of the coupling y_1 is plotted against the DM mass. The $\psi\psi \rightarrow S^+ S^-$ annihilation channel opens up at different S^+ values, as shown in the left plot in Fig. 6.4.1, for a fixed χ^+ mass of 1 TeV. In the right plot of Fig. 6.4.1, m_χ is varied for the entire range of m_ψ keeping the mass splitting fixed at two specific values. For smaller mass splitting, a relatively lighter propagator reduces the demand on the requirement of large coupling. However, to be compatible with the direct detection limits of $y_1 \leq 0.05$, a highly degenerate region of the parameter space may be required, as explained towards the end of the next subsection. It is to be noted here that if S^+ was not charged, the one-loop penguin diagrams (Fig. 6.D.1) contributing to direct search would not contribute, but still, the strength of the new physics coupling would determine the interplay between the nonthermal and thermal production, even leading to freezeout for a wider parameter space. Because of this generic feature, we

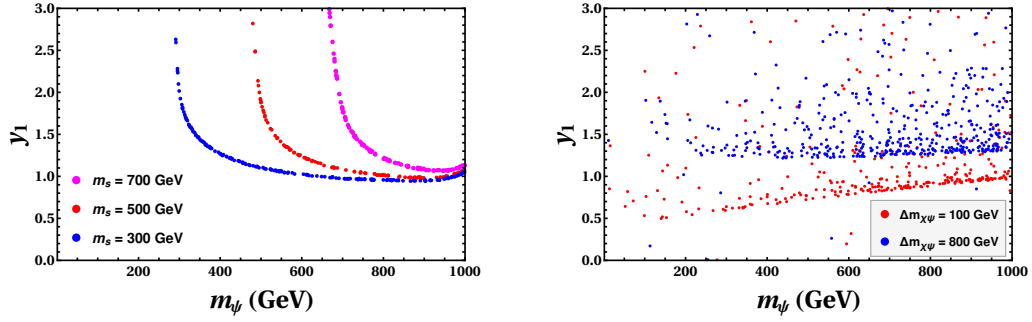


Figure 6.4.1: y_1 vs m_ψ for specific values of m_S in agreement with the measured DM relic density; Left: for a fixed propagator mass of $m_\chi = 1$ TeV; Right: for varying m_χ with the mass splitting fixed at $m_\chi - m_\psi = 100$ GeV (red) and 800 GeV (blue) keeping $65 \leq m_S \leq 1000$ GeV.

have kept the full range of y_1 in the analysis, although for our particular case freezeout is restricted to a narrow window of the parameter space.

6.4.2 Region with almost degenerate case: $m_\chi \sim m_\psi$

In this region, as the mass difference between χ^+ and ψ is very small, the effect of the co-annihilation channels become important. Unlike Fig. 6.3.2(a), these channels provide correct relic density even when $m_\psi \leq m_S$. This is because ψ can now annihilate into S^+ and a photon or a Z boson provided $2m_\psi \geq m_S/(m_S + m_Z)$ (Fig. 6.3.2(c)-(d)). Setting a non-zero value of λ_1 enables the process considered in the diagram in Fig. 6.3.2(b), although the effect of this cross section is very small. However, since λ_1 is fixed at a small value ($= 0.01$) in the analysis, this contribution is always negligible. We consider the ranges of parameters given in Table 6.4.2 while performing the scan.

$1 \text{ GeV} \leq m_\psi \leq 1 \text{ TeV}$	$m_\chi = m_\psi + 1, 5, 10 \text{ GeV}$
$65 \text{ GeV} \leq m_S \leq 1 \text{ TeV}$	$0.0 \leq y_1 \leq 3.0$

Table 6.4.2: Range of parameters considered for the scan in Region 2: $m_\chi \sim m_\psi$ with thermal freeze-out of dark matter ψ .

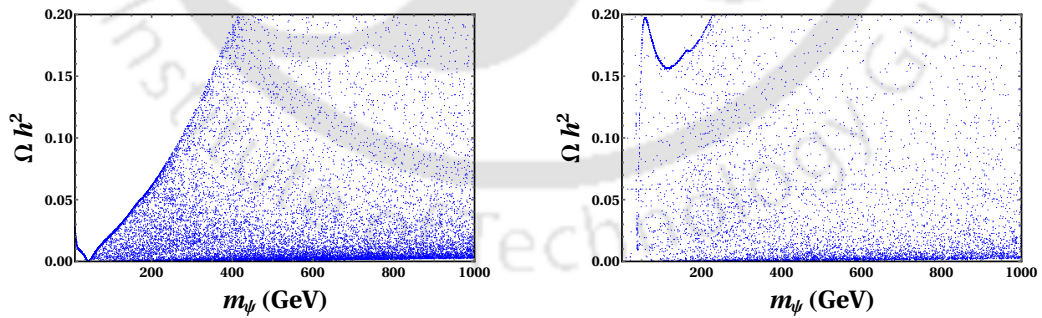


Figure 6.4.2: DM relic density vs DM mass for mass splitting between dark sector particles fixed at 1 GeV (left) and 5 GeV (right). For larger mass splitting (right), the effect of Z -mediated s -channel cross sections arising due to χ^+ annihilation become negligible, and co-annihilation channels dominate increasing the relic density in the low DM mass regime within observed limits.

In Fig. 6.4.2, the variation of DM relic density is plotted with DM mass for two different mass splittings between χ^+ and ψ . In the left panel, the splitting is fixed at 1 GeV. For such a small splitting, the annihilation of the other dark sector particle χ^+ into SM fermions also contributes to the relic density along with the diagrams given in Fig. 6.3.2. All these additional channels being Z -mediated, we see the Z -resonance at $m_\psi = 45$ GeV. After the resonance, these cross sections decrease with an increase in m_ψ making the co-annihilation channels important. But if the mass splitting is made larger (5 GeV as

shown in the right plot), then the effect of mediator annihilation, ie, the contribution of these Z -mediated diagrams becomes negligible compared to ψ co-annihilation and pair annihilation, so that even the low mass regime of m_ψ supports the correct relic density. In Fig. 6.4.3, the variation of y_1 vs m_ψ is plotted

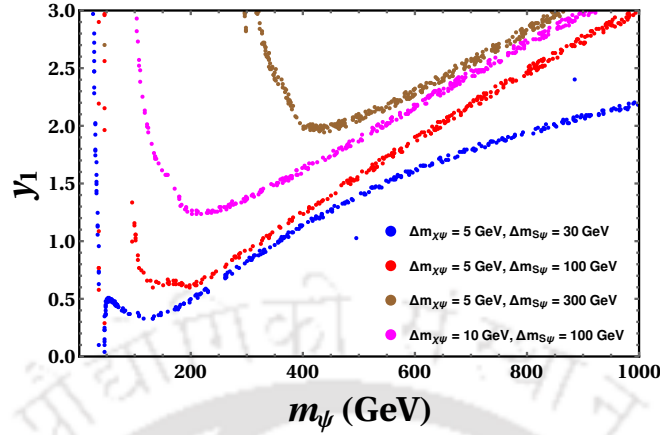


Figure 6.4.3: y_1 vs m_ψ plot indicating the significance of the co-annihilation channels, with the mass splitting $\Delta m_{\chi\psi} = m_\chi - m_\psi = 5$ and 10 GeV, and for specific values of $\Delta m_{S\psi} = m_S - m_\psi$.

with the mass splitting between ψ and χ^+ ($\Delta m_{\chi\psi}$) and the splitting between χ^+ and S^+ ($\Delta m_{S\psi}$) fixed at specific values. The dip at $m_\psi = 45$ GeV indicates a broad s -channel Z -resonance arising due to $\chi^+\chi^-$ mediator annihilation into SM fermions. The corresponding cross section being proportional to $e^{-2\Delta m_{\chi\psi}/m_\psi}$, it becomes irrelevant at the mass splitting of 10 GeV, as indicated by the absence of such a dip in the deep pink curve. The co-annihilation channels, on the other hand, has a slightly more complex dependence on the masses. Considering the Feynman diagrams given in Fig. 6.3.2, three channels could contribute. Two s -channel processes mediated by S^+ and a t -channel process mediated by χ^+ . Both the s -channel processes have S^+ in the final state as well. This brings in a tug of war between the resonant condition, $m_\chi + m_\psi = m_S$ and the favourable phase space for the lighter final state particles, consequently disfavouring the $\chi^+\psi \rightarrow S^+h$ channel. On the other hand, the other s -channel process with $S^+\gamma$ in the final state becomes significant around $m_S = m_\chi + m_\psi \sim 2m_\psi$. The t -channel process takes over at larger m_ψ , and the cross section slowly saturates. This is indicated in the plots at higher ψ masses. Heavier S^+ requires larger couplings to compensate for the diminishing effects of the propagator and narrowing phase space. This is visible in the plots with larger $\Delta m_{S\psi}$ values. As soon as m_ψ becomes larger than m_S , the pair annihilation of ψ becomes dominant and the correct relic density is observed for m_ψ up to the TeV scale. Summing up, the solutions compatible with the required $y_1 \lesssim 10^{-5}$ is possible with highly degenerate regions with $\Delta m_{\psi\chi} \lesssim 5$ GeV in a narrow mass range of $m_\psi \sim m_Z/2$.

6.5 Collider signatures

The signatures of the model can be traced in colliders like the LHC through the gauge production and subsequent decay of the charged partner fermion, χ^+ and the charged scalar S^+ . We plot the cross section for pair production of χ^+ (σ_χ) at the 14 TeV LHC in Fig. 6.5.1a). At $m_\chi = 200$ GeV, the cross section is of the order of 100 fb, which falls down to 0.5 fb for $m_\chi = 1$ TeV. The charged scalar S^+ , on the other hand, has much smaller cross section (σ_S) of about 1 fb for $m_S = 100$ GeV, which reduces to about 0.1 fb for $m_S = 500$ GeV as seen from Fig. 6.5.1b). Subprocesses contributing to σ_χ and σ_S are $pp \rightarrow \chi^+\chi^-/S^+S^-$ via Z/γ s -channel diagrams. The decays of χ^+ and S^+ are controlled by the Yukawa couplings y_1 and y_2 . When both the couplings are large enough, the decays are instantaneous and happens at the interaction point itself where the particles are produced. On the other hand, for small values of the couplings, the decay width could be considerably smaller so that the decays happen away from the interaction point, leaving distinct signatures characteristic to such Long-Lived Particles

(LLP). Indeed, such LLP arising in different contexts are actively studied for their signatures at LHC and other exclusively designed colliders like MATHUSLA in the past few years. Signatures typically include (i) *Displaced Vertex*, where the LLP decay within the detector, but away from the interaction point; (ii) *Heavy Stable Charged Particle (HSCP) tracks*, where charged LLP decaying outside the detector, leaving charged tracks similar to that of the muon, but with distinct features arising from slower speed and larger energy deposition along with its flight through the materials of the detector or (iii) *Disappearing Tracks*, where an isolated track stopping before hitting the outer layers of the silicon tracker and without any associated energy deposition in the calorimeters or muon chamber.

Experimental searches by CMS [211, 213, 214], ATLAS [4, 6–8] and LHCb [10, 11, 187] considered long-lived neutral scalar decaying in the calorimeter into pair of jets, looking for displaced vertex with negative results leading to constraints on the masses of the corresponding LLP. Considering a gauge-mediated supersymmetry breaking scenario, CMS [212] considered non-prompt jets arising from the decay of gluino to gluons and gravitinos, where the gluon jet emerges from a displaced vertex within the tracker. HSCP search by ATLAS [2, 3] and CMS [178] looking for possible candidates in the supersymmetric scenarios have obtained mass limits on the so-called R -hadrons (long-lived gluinos or squarks), staus and charginos. A long time of flight along with the presence of anomalously high energy deposits in the silicon detector is considered as the signature of the HSCP. Ref. [2] has used the ionisation information from the pixel subsystems of the ATLAS detector to probe the HSCP modelled after R -hadron. CMS [107, 210] and ATLAS [5] considered disappearing tracks arising from chargino decay to neutralinos and a low-momentum pion, where the neutralino is nearly mass degenerate with the chargino, resulting in its long decay time. A recent report on the strategies that may be adopted in searching for LLP at colliders like the LHC discussing simplified models, experimental signatures, details of the backgrounds and possible requirements of detector upgrades presented in Ref. [22]. A comprehensive analysis of LLP arising in FIMP like scenario is presented in Ref. [49], where the authors have described possible signatures of a long-lived charged fermion decaying into the SM charged lepton and scalar DM candidate. Considering fermionic dark matter as a mixture of $SU(2)_L$ singlet and doublet LLP signatures at the LHC are studied in Ref. [84]. Charged dark doublet scalar decaying into fermionic FIMP and charged leptons are considered in [162]. In addition to the possible searchers at the LHC, there are proposals for dedicated detectors like MATHUSLA [114, 191], which can potentially look for long-lived neutral particles decaying to two jets or charged leptons outside the LHC detectors. These surface detectors are located to look for long-lived neutral particles originating from the interaction point, decaying in its volume into two identifiable jets or leptons.

Many of the signatures discussed above would be present in the scenario discussed in this chapter. In this case, the charged dark fermion, χ^+ , the $\mathbb{Z}_2$ -even charged scalar, S^+ and the heavy neutrino N are all possible candidates of LLP. The specific signatures of the LLP will depend on many aspects. A particular feature about the model presented in our case is the possibility of *multiple LLP signatures* arising through the cascade decay of the kind $LLP1 \rightarrow LLP2 + X$. In the present case, this is possible whenever the 2-body decay of χ^+ is allowed. In such case, χ^+ decays to an on-shell S^+ , which itself could be an LLP when its decay is slow, enabled by small y_2 . However, for large enough y_2 , S^+ decays quickly enough without leaving any displaced vertex. In that case, the LLP signature of χ^+ is the standard case of usual single LLP. Interesting signature of multiple LLP processes is not investigated in detail in the literature. To illustrate this we shall take a specific case with $c\tau(\chi) = \frac{c}{\Gamma_\chi} \leq 50$ cm and $c\tau(S^+) \sim 1 - 10$ m, with the first displaced vertex falling in the Inner Detector, followed by the second one (that of S^+ decay) falling in the Calorimeter or Muon Spectrometer. The first displaced vertex will leave a kink on the track, with the accompanying ψ invisible. Distinguishability of this kink depends on the masses and energy of the particles, apart from the granularity of the detectors. While we recognise the fact that such possibilities are complex to probe, with a precision of a few microns on the position in the pixel detector, and high precision calorimeters it is worth exploring the possibilities. In addition to charged χ^+ and S^+ , we have the presence of heavy neutral fermions that couple to the charged leptons through the standard Higgs boson. The Yukawa coupling corresponding to this interaction is small, making N almost always stable within the detector. However, for N heavier than 125 GeV, Yukawa

couplings of the order of 10^{-8} would lead to its decay within the detector.

We shall consider some preliminary numerical analysis of the above long-lived situations with χ^+ , S^+ and N separately in the following subsections.

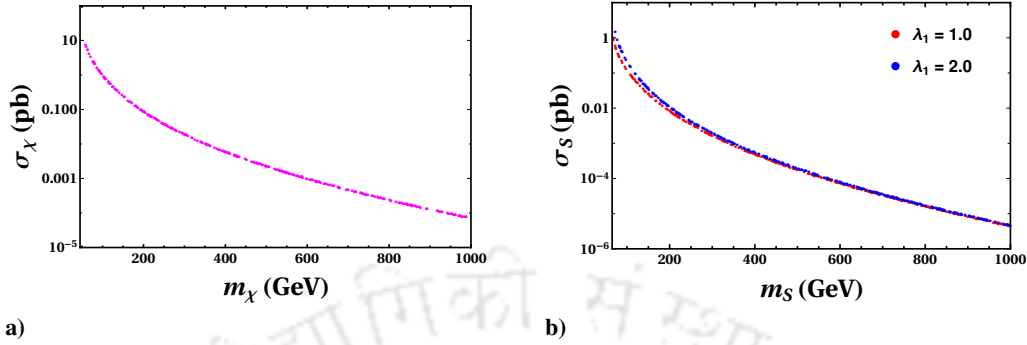


Figure 6.5.1: Production cross section of $\chi^+\chi^-$ (a) and S^+S^- (b) pair at 14 TeV LHC for different values m_χ and m_S , respectively.

6.5.1 Long-lived charged fermion, χ^+

The conditions on χ^+ arising from the dark matter considerations depend on its mass and nature of decay. In the kinematic regions where $m_\chi > m_S + m_\psi$, the decay of χ^+ is dictated by the Yukawa coupling, y_1 , the Feynman diagram for which is given in Fig. 6.3.1a). From Fig. 6.3.3 it is clear that χ^+ lighter than about 200 GeV is not favoured in this case. The compatible couplings are restricted to $10^{-13} \leq y_1 \leq 10^{-8}$, leading to $c\tau$ in the range of interest of LHC for the larger limits of the coupling. Fig. 6.5.2a) gives the decay lengths for a range of m_χ and typical values of the coupling, for fixed masses of $m_\psi = 1$ GeV and $m_S = 65$ GeV. We have checked the sensitivity with varying m_ψ and m_S and found that the decay length remains in the same order of magnitude for a large range of these masses. Keeping all other parameters the same, the dependence of $c\tau$ on m_χ is minimal, allowing a wide range of the mass values to be accessible as LLPs at the LHC. We tabulate the decay lengths for typical values of m_χ fixing $y_1 = 10^{-8}$ in Table 6.5.1. Note that the decay width is proportional to y_1^2 in this case, and therefore, $c\tau$ for a specific y_1 value can be obtained from what is quoted in a table by using the scaling factor $10^{-16}/y_1^2$. For example, by choosing a slightly larger y_1 value the decay length can be brought down to the centimetre range, which is relevant to the Inner Detector of LHC experiments. At the same time, couplings smaller than 10^{-8} leads to χ^+ escaping the detector without decaying, leaving charged track in all sections of the detector, very similar to that of the μ tracks, but with larger energy deposition owing to it being heavier [49, 103]. Besides, its time of flight will be larger than that of the muon as it would be slower than muons.

Moving on to the case of $m_\chi < m_S + m_\psi$, the two-body decay is kinematically disallowed and the three-body decay $\chi^+ \rightarrow \psi N \ell^+$ is controlled by the product $y_1 y_2$ and mediated by a virtual S^* , as shown in the Feynman diagram in Fig. 6.3.1b). In this case, χ^+ can be lighter than 100 GeV as seen from Fig. 6.3.4, with the product of the couplings mostly range between 10^{-10} and 10^{-6} . With conditions of mass splitting satisfied for enhanced co-annihilation, possibilities with larger couplings even up to 10^{-5} are not rare, and may even reach up to 10^{-4} albeit with less probability, as presented in Table 6.3.3. In Fig. 6.5.2b) the decay length $c\tau$ is plotted against m_χ for exemplary values of the relevant coupling combination, keeping $m_\psi + m_N + m_\ell < m_\chi < m_\psi + m_S$ so as to focus on the 3-body decay. The dependence on the mass here is stronger than that in the 2-body decay case discussed above. For the entire mass range of $m_\chi = 40$ GeV to large values of $m_\chi = 1000$ GeV, it is possible to have χ^+ decaying within the Inner Detector or the Calorimeters, or even decay outside the detectors, leaving an LLP signature. Table 6.5.1(b) shows the decay lengths for some specific benchmark points relevant to the LHC.

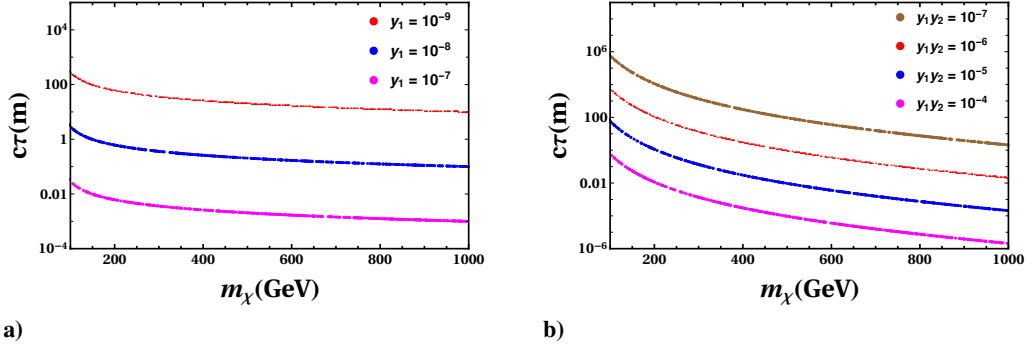


Figure 6.5.2: Decay length vs mass of χ^+ for different couplings corresponding to (a) 2-body and (b) 3-body decay of χ^+ . The neutral particle masses are $m_\psi = 1$ GeV and $m_N = 30$ GeV, while the charged scalar masses m_S are 65 GeV for case (a) and 1 TeV for case (b).

y_1	m_χ	m_ψ	m_S	$c\tau$
	[GeV]			[m]
10^{-8}	200	130	65	55
		65	65	72
		65	130	110
	1000	900	65	63
		65	65	88
		65	900	710

(a)

$y_1 y_2$	m_χ	m_ψ	m_S	$c\tau$
	[GeV]			[m]
10^{-6}	200	130	65	55
		65	65	72
		65	130	110
	1000	900	65	63
		65	65	88
		65	900	710

(b)

Table 6.5.1: Decay length ($c\tau$) of χ^+ in meters for typical masses of ψ and S^+ and at a fixed value of the coupling. Heavy neutrino mass is set to $m_N = 30$ GeV.

6.5.2 Long-lived charged scalar, S^+

Turning to the case of S^+ , its decay length is decided purely by the strength of y_2 . For the freeze-in solution of dark matter, we needed $y_1 \lesssim 10^{-8}$ to be compatible with the observed relic density. The charged scalar, S^+ is mostly produced through gauge mediation, with a small contribution from the Higgs mediated process enabled by the quartic coupling λ_1 . When the coupling y_1 is of the order of unity, on the other hand, as required in the case of freeze-out mechanism to generate the dark matter relic density, χ^+ would decay instantly as it is produced. It will decay to $S^+ \psi$, with S^+ further decaying and ψ missing. The cross section for this is given in Fig. 6.5.1b) against m_S at 14 TeV LHC. The decay of S^+ is decided by the Yukawa coupling y_2 . The decay width is given by

$$\Gamma_{S^+ \rightarrow Nl} = \frac{y_2^2}{16\pi m_S^3} (m_S^2 - m_N^2 - m_l^2) [m_S^2 - (m_N + m_l)^2]^{1/2} [m_S^2 - (m_N - m_l)^2]^{1/2}. \quad (6.5.1)$$

In Fig. 6.5.3 we plot the decay length of S^+ against m_S for specific choices of the relevant Yukawa coupling, y_2 . Considering this to be of the order of $10^{-8} - 10^{-10}$ makes it decay within the LHC detector, but with a measurable delay from its production to distinguish it from other particles which decay instantly. The signatures depend on the actual decay width. For $y_2 \sim 10^{-8} - 10^{-9}$ it leave a displaced vertex in the Inner Detector, whereas weaker couplings would make it decay within the Calorimeters or even in the muon detectors. Beyond $y_1 \sim 10^{-10}$ will see S^+ escaping the LHC detectors

without decaying, but leaving a clear charged particle track in all sectors. Similar to the case of χ^+ discussed above, this can perhaps be distinguished from the muon tracks, owing to its heavy nature and slower speed. It is also interesting that the mass splitting between S^+ and N dictates the signatures of S^+ at the collider. If the splitting is small, a soft lepton will be produced in the final state, leaving disappearing tracks. On the other hand, if the splitting is large, then very likely a kink will be seen in the typically charged track signal. The cross-sections and decay widths of S^+ for selected benchmark points are given in Table 6.5.2 for some typical benchmark points which satisfy the relic density bound.

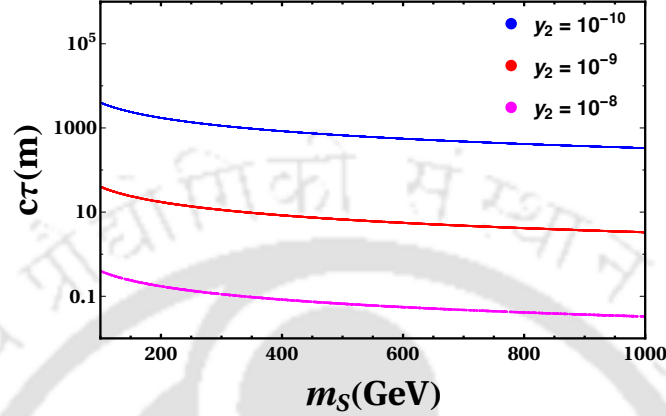


Figure 6.5.3: Decay length vs mass of long lived particle (S^+) for different couplings. m_N is fixed at 30 GeV.

	m_ψ	m_S	m_χ	y_1	Ωh^2	$\sigma_{pp \rightarrow S^+ S^-}$ (pb)	$\Gamma_{S^+ \rightarrow N \ell}$ (GeV)	$c\tau$ (m)
BP 1	56	165	473	7.53×10^{-13}	0.12	0.0144	9.21×10^{-18}	21.5
BP 2	101	77	613	1.18×10^{-13}	0.118	0.2441	3.3×10^{-18}	60.0
BP 3	225	188	228	0.01	0.12	0.0088	1.07×10^{-17}	18.5
BP 4	150	127	154	0.05	0.119	0.0376	6.76×10^{-18}	29.5

Table 6.5.2: Values of benchmark points considered in this paper for the production cross section (at 14 TeV LHC) and decay of the charged scalar S^+ . DM relic density is achieved through either freeze-in (2 body) or freeze-out. All masses are in GeV. For all points, $m_N = 30$ GeV and $y_2 = 10^{-9}$.

6.5.3 Long-lived neutral fermion, N

The neutral heavy fermion, N , decays through $N\nu h$ interaction, where ν is the SM neutrino and h is the SM Higgs boson. The strength of this interaction is decided by the Yukawa coupling y_N , which also decides the neutrino mass through the seesaw mechanism. As discussed in the introduction, this coupling is required to be $\sim 10^{-8}$ for m_N considered in our discussion. Also, for $m_N < 125$ GeV it cannot decay to an on-shell Higgs boson, further slowing down the decay. Therefore, N (of mass in the 100 GeV range or lighter) decays much beyond the LHC detectors, consequently leaving missing energy signature, like the SM neutrinos. Dedicated detectors like the MATHUSLA [114] placed beyond the LHC detector complex could capture the signatures of such long-lived particles. This surface detector is typically prepared to identify decay of LLP into a pair of charged particles, which are detected by its multi-layer trackers. However, notice that N decays to neutrino and Higgs boson. With the missing energy, it will be difficult to correlate the event as originating from the interaction point. Different techniques like time of flight information may need to be made use of, requiring close talks between the CMS/ATLAS and MATHUSLA detectors. On the other hand, for $m_N \gtrsim 130$ GeV, N can have decay length within the CMS/ATLAS detector, and can in principle be probed through the displaced vertex

study. Again, the typical search strategies adopted in the case of displaced vertex studies of ATLAS and CMS will not work in this case. It may be required to look for reconstructed Higgs (produced almost at rest, especially when N is nearly degenerate with Higgs) originating from a displaced vertex along with associated missing transverse momentum to the direction of the interaction point could be a possible signature. Although we have not considered heavier N for our analysis, to get the same relic density,

m_N (GeV)	Γ_N (GeV)	$c\tau$ (m)
100	1.53×10^{-22}	1.3×10^6
120	7.53×10^{-22}	2.8×10^5
128	4.09×10^{-19}	480.5
130	1.1×10^{-18}	179
135	4.1×10^{-18}	48.05
150	2.09×10^{-17}	9.42
200	1.1×10^{-16}	1.79

Table 6.5.3: Decay width and decay length of N . $y_N = 10^{-8}$ in all cases. It is to be noted that the order of Γ_N changes substantially for $m_N \geq 125$ GeV, which implies the onshell production of SM Higgs.

the increase in m_N in the decay width of χ can be compensated by larger $y_1 y_2$.

Finally, we may consider these search options in the proposed e^+e^- colliders like CLIC with sufficient centre of mass energy to produce the charged particle, χ^+ and S^+ . The search strategies in this case would be very similar to those mentioned above in the case of the LHC. Notice that the production being electroweak process, the cross section, in this case, will be comparable with the LHC. On the other hand, the clean environment and fixed centre of mass energy can be of great advantage in the leptonic colliders. We shall defer a detailed study to a future publication.

6.6 Summary

Providing a viable dark matter candidate, yet compatible with all the experimental observations is one of the major challenges of particle physics today. In the popular WIMP scenario, weakly interacting dark matter particles produced thermally in abundance is depleted with large enough annihilations to achieve the observed relic density as they decouple and freeze out. In most models of WIMP, the basic interaction that leads to such annihilations is the same as that is relevant in the direct detection experiments based on their elastic scattering from nuclei. Presently, the direct detection experiments have constrained these couplings to the extent that it becomes quite difficult to accommodate the required annihilation cross section. This has led many researchers to favour an alternate mechanism, whereby very feebly interacting massive particles are slowly produced to build up the relic density of the dark matter. In this FIMP scenario, usually, a thermally produced partner particle decays very slowly to the dark matter particle.

In this chapter we consider a gauge singlet fermionic FIMP dark matter, ψ and an isospin singlet charged fermion, χ^+ as its partner particle. The stability of the dark matter particle is ensured by imposing a $\mathbb{Z}_2$ symmetry under which both χ and ψ are odd, while all other particles are even. Gauge produced χ^+ decays through its Yukawa coupling (y_1) to ψ and another newly introduced isospin singlet charged scalar S^+ . For very small values of this Yukawa coupling, the dark matter density will be slowly built up. Our analysis shows that for $10^{-13} \leq y_1 \leq 10^{-8}$ a wide range of DM mass 1 GeV $\leq m_\psi \leq 600$ GeV is viable for partner mass of 100 - 1000 GeV. The presence of S^+ and the Yukawa coupling (y_2) enabling its decay to a heavy neutral fermion and the charged SM leptons, add an entirely

new dimension to the dynamics compared to what is usually discussed in the literature. The signatures here include a slowly decaying long-lived particle, as well as the dark matter particle itself. Under favourable conditions, the partner particles can be produced in sufficient amounts at colliders like the LHC, which then will decay away from the interaction point, but within the detector, leaving distinct signatures compared to many exotic particle searches. The typical coupling (relevant to the decay) required for this is $\sim 10^{-9}$, which in the usual FIMP scenarios discussed in the literature require very light dark matter candidates (keV - MeV). In the present scenario we discussed, the presence of y_2 , which can be independent of the dark matter considerations makes it possible to have both χ^+ and S^+ to be a long-lived particle that could be searched at the LHC, for a large range of its mass. When the kinematic condition $m_\chi > m_\psi + m_S$ is not met, χ^+ decays through off-shell S^+ to $N\ell^+\psi$. This decay is controlled by a combination of y_1 and y_2 . With y_1 within a restricted range, y_2 can be arranged so that the delayed decay of χ^+ to the DM candidate ψ adds significantly to its number density. This is a new feature of the dark matter dynamics that is comparatively less explored in the literature. This intermingling of slow production leading to freeze-in and annihilation followed by freeze-out opens up a large region of the parameter space, otherwise unavailable when these mechanisms are considered separately. We have discussed the details of this in Section 6.3.2, where the importance of the annihilation process is established. In sharp contrast to the usual freeze-in plus LLP scenario, the DM mass range that is compatible in the present case is from a few GeV to TeV. While the usual HSCP scenario arises in certain parameter regions, other distinct features are requiring modified search strategy at LHC/MATHUSLA detectors as summarised below.

1. Having two charged *LLP*'s, it is possible to have $LLP1 \rightarrow LLP2 + \cancel{E}$, with *LLP2* decaying outside the detector. With large enough missing energy, this can lead to an HSCP track with a kink at the first decay vertex.
2. $LLP1 \rightarrow LLP2 + MET \rightarrow SM + MET$. Within favourable parameter regions, this can lead to sufficiently large disappearing track, with a kink at the first decay vertex.
3. The neutral *LLP* decay to Higgs boson and a neutrino ($N \rightarrow h\nu$). With h further decaying to $b\bar{b}$, this could lead to two jets with displaced vertex, and having invariant mass around Higgs mass, within ATLAS/CMS or MATHUSLA. The reconstructed Higgs momentum will not point to the interaction point owing to large associated missing energy of the undetected neutrino.

We have explored various possibilities, and show that for m_ψ starting from 1 GeV up to 1000 GeV and beyond, almost the entire range of theoretically allowed coupling values and a large range of masses of other particles are compatible with the observations. The collider phenomenology in this region of parameter space may not feature any long-lived particle. On the other hand, the presence of two exotic charged particles, one fermion and one scalar will add rich phenomenology, which will be explored in detail in a future study. For easy reference, we have summarised the conclusions in a qualitative way for different ranges of the relevant couplings in Table 6.6.1.

Finally, we point the attention of the reader to the presence of heavy neutral fermion inducing type-I seesaw mechanism providing a way to generate light neutrino mass within the same framework, as an added feature of the model.

Appendices

6.A Boltzmann Equations

The dark matter particle ψ when produced with the slow decay of χ^+ , but with a rate that results in the sufficiently large number density of ψ before freezing-in presents a novel feature as discussed towards the end of Section 6.3.2. With a sufficiently large value of the relevant coupling, y_1 , the thermal production and pair annihilation become significant, leading to freeze-out. The two limiting cases are (i)

DM scenario	couplings			LLP	collider signature
	y_1	y_2	$y_1 y_2$		
Freeze-in via 2 body decay	$10^{-13} \lesssim y_1 \lesssim 10^{-9}$	$10^{-10} \lesssim y_2 \lesssim 10^{-7}$	small	χ^+	charged track
				S^+	displaced vertex, charged track
Freeze-in via 3 body decay	$\lesssim 10^{-8}$	$10^{-3} \lesssim y_2 \lesssim 1$	$10^{-11} \lesssim y_1 y_2 \lesssim 10^{-8}$	χ^+	charged track
Thermal + non-thermal production via χ^+ decay (dominant)	$10^{-8} \lesssim y_1 \lesssim 10^{-7}$	$10^{-3} \lesssim y_2 \lesssim 1$	$10^{-11} \lesssim y_1 y_2 \lesssim 10^{-7}$	χ^+	charged track
Freeze-out	~ 1	$\lesssim 10^{-7}$	$\lesssim 10^{-4}$	χ^+	charged track, displaced vertex
				S^+	displaced vertex, charged track
Freeze-out	~ 1	~ 1	~ 1	—	—

Table 6.6.1: Table showing the possible DM scenarios and the limits on the coupling from relic density and collider requirements. LLP possibilities and signatures are also mentioned.

with y_1 is small enough so that only the production through χ^+ decay is relevant, and (ii) with y_1 large enough so that the thermal production and subsequent freeze-out scenario are relevant. The Boltzmann Equation controlling the yield, in this case, depends on the annihilation cross section of ψ along with the decay rate of χ^+ as given by

$$\frac{dY_\psi}{dx} = \sqrt{\frac{45}{4\pi^3 G}} \frac{x}{m_\psi^2} \frac{\sqrt{g_*(T)}}{g_s(T)} \langle \Gamma_{\chi \rightarrow \psi X} \rangle Y_{\chi_{eq}} - \sqrt{\frac{\pi}{4\pi G}} \frac{m_\psi}{x^2} \sqrt{g_*(T)} \langle \sigma v \rangle_\chi (Y_\psi^2 - Y_{\psi_{eq}}^2),$$

with $Y_{\chi_{eq}}$ and $Y_{\psi_{eq}}$ are the equilibrium number densities for χ^+ and ψ respectively, $x = m_\psi/T$, where T is the photon temperature. Notice that in the standard freeze-in case x is defined in terms of the mass of the decaying particle (χ^+ in our case). However, what is relevant in the present case is the final freeze-out, and therefore we have considered $x = m_\psi/T$, appropriately redefining other factors.

$$\sqrt{g_*(T)} \simeq \frac{g_s(T)}{\sqrt{g_\rho(T)}}$$

where g_s and g_ρ are the degrees of freedom corresponding to entropy and energy density of the Universe.

$$\langle \Gamma_{\chi^+ \rightarrow \psi X} \rangle = \Gamma_{\chi^+ \rightarrow \psi X} \frac{K_1(x')}{K_2(x')}$$

$x' = m_\chi/T$, K_1 and K_2 are modified Bessel functions and $\Gamma_{\chi^+ \rightarrow \psi X}$ represents the decay width of χ^+ to ψ and X , where X denotes S^+ in case of 2 body and N and l^+ for 3 body decay of χ^+ . $\langle \sigma v \rangle_\chi$ denotes the thermal average of all possible pair annihilation and coannihilation cross-sections of ψ .

The relic density is related to the yield through the standard relation,

$$\Omega_\psi h^2 = 2.75 \times 10^8 m_\psi Y_\psi(T_0), \quad (6.A.1)$$

where T_0 is the present temperature of the universe.

6.B Constraint from $H \rightarrow \gamma\gamma$

The Higgs boson has a direct quartic coupling with the new charged scalar, S^+ ,

$$\mathcal{L} \ni \lambda_1 S^\dagger S H^\dagger H. \quad (6.B.1)$$

The trilinear coupling arising from this after the electroweak symmetry breaking results in S^+ contribution to $H \rightarrow \gamma\gamma$ at the leading order through the triangle diagram. Including this contribution from S^+ , the diphoton decay width is given by [123, 149]

$$\Gamma(H \rightarrow \gamma\gamma) = \frac{\alpha^2 m_H^3}{256 \pi^3 v^2} \left| F_W + \sum_f N_{cf} Q_f^2 F_f + \lambda_1 v F_S \right|^2, \quad (6.B.2)$$

where

$$F_W = 2 + 3\tau_W [1 + (2 - \tau_W) f(\tau_W)], \quad (6.B.3)$$

$$F_f = -2\tau_f [1 + (1 - \tau_f) f(\tau_f)], \quad (6.B.4)$$

$$F_S = \tau_S [1 - \tau_S f(\tau_S)], \quad (6.B.5)$$

with $\tau_i = 4m_i^2/m_h^2$ and

$$f(\tau) = \begin{cases} [\sin^{-1}(1/\sqrt{\tau})]^2 & \text{for } \tau \geq 1, \\ -\frac{1}{4} \ln \left(\frac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}} \right) - i\pi & \text{for } \tau < 1. \end{cases} \quad (6.B.6)$$

With $m_S = 150$ GeV, we have

$$\Gamma(H \rightarrow \gamma\gamma) \sim 0.25 \text{ (keV)} (6.54 - 90.5 \lambda_1)^2. \quad (6.B.7)$$

6.C Constraint from Invisible Higgs decay

Higgs boson does not have a direct coupling to the dark matter, ψ . However, $H \rightarrow \psi\psi$ is possible

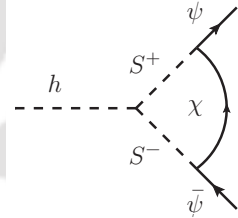


Figure 6.C.1: Feynman diagram contributing to the invisible Higgs decay through $H \rightarrow \psi\psi$.

through the one-loop diagram as in Fig. 6.C.1. This contribution to the invisible decay of the Higgs is given by [159]

$$\Gamma(H \rightarrow \psi\psi) = \frac{m_H}{2\pi} \left[1 - \frac{4m_\psi^2}{m_H^2} \right]^{3/2} |g_V|^2, \quad (6.C.1)$$

with

$$g_V = \frac{\lambda_1 y_1^2 v}{32 \pi^2} \left[2 m_\psi C_1(m_\psi^2, m_H^2, m_\psi^2, m_\chi, m_S, m_S) - m_\chi C_0(m_\psi^2, m_H^2, m_\psi^2, m_\chi, m_S, m_S) \right],$$

where C_i are the relevant Passarino-Veltman loop functions. With all new particle masses around GeV-TeV scale, $\Gamma(H \rightarrow \psi\bar{\psi}) \sim \mathcal{O}(10^{-6})$ for the couplings $y_1 \sim 1$ and $\lambda_1 \sim 0.01$ which is negligible compared to the standard contributions.

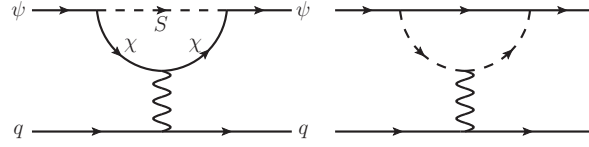


Figure 6.D.1: One-loop diagrams that contribute to dark matter - nucleon scattering. Notice that there are no tree-level contributions.

6.D Direct Detection Constraints

There are two sets of Feynman diagrams with photon and Z mediation, as in Fig. 6.D.1. The spin-independent direct detection cross section is given by [94]

$$\sigma_{\text{SI}}^{\psi\chi} = \frac{16}{\pi} \frac{m_\psi^3 m_p^3}{m_\psi + m_p} |c_{VV}^q|^2, \quad (6.D.1)$$

where m_p is the proton mass. The relevant Wilson coefficients are given by

$$\begin{aligned} c_{VV}^q = & -\frac{\alpha_{\text{em}}}{24\pi x_\psi^4} \frac{1}{m_S^2} Q Q_q y_1^2 \\ & \times \left\{ \left[-3x_\psi^6 + 6x_\psi^5 x_\chi + 12x_\psi x_\chi (1 - x_\chi^2)^2 + 8(1 - x_\chi^2)^3 + 2x_\psi^4 (5 + x_\chi^2) \right. \right. \\ & \left. \left. - 6x_\psi^3 x_\chi (1 + 3x_\chi^2) - 3x_\psi^2 (5 - 2x_\chi^2 - 3x_\chi^4) \right] \frac{g(x_\psi, x_\chi)}{1 - (x_\psi - x_\chi)^2} \right. \\ & \left. + \frac{2x_\psi^2 (4 - 3x_\psi^2 + 6x_\psi x_\chi - 4x_\chi^2)}{1 - (x_\psi - x_\chi)^2} + (8 + x_\psi^2 - 4x_\psi x_\chi - 8x_\chi^2) \ln x_\chi \right\}. \end{aligned} \quad (6.D.2)$$

The loop function in terms of the mass ratios $x_\psi \equiv m_\psi/m_S$ and $x_\chi \equiv m_\chi/m_S$ is given by

$$g(x_\psi, x_\chi) = \frac{\ln \left(1 - x_\psi^2 + x_\chi^2 + \sqrt{x_\psi^4 + (1 - x_\chi^2)^2 - 2x_\psi^2(1 + x_\chi^2)} \right) - \ln(2x_\chi)}{\sqrt{x_\psi^4 + (1 - x_\chi^2)^2 - 2x_\psi^2(1 + x_\chi^2)}}. \quad (6.D.3)$$

Conclusion

This doctoral thesis focuses on developing theoretical models providing particle dark matter candidates, as well as studying the possibilities of testing them at high energy particle colliders, such as the Large Hadron Collider (LHC). As we elaborate in the chapter outlines below, we have worked on several scenarios which efficiently address some of the shortcomings of single WIMP DM candidate, and at the same time, has intriguing implications in the DM dynamics and collider probes. As a brief outline, we have primarily addressed three such scenarios beyond the single-WIMP stereotype :

Part I: Multi-component DM : In this part, we take up the possibility of multi-component DM scenario [56, 65] as a viable possibility where the current experimental bounds eliminated the parameter space of the conventional WIMP paradigm. Instead of a single DM candidate, as is the case for the conventional WIMPs, here we consider two WIMP DM candidates with possibilities of interaction between them. The interplay between the two DM candidates here is the key point of the works presented here [89, 93]. This is because this interaction is entirely independent of the coupling with the visible sector, and thus, through this coupling, relic density can be successfully explained despite keeping direct detection bounds below the observed limit. Here, we discussed two works on DM components of different spin and isospin in this context, with the indication of possibilities of collider probes.

Chapter 2 : First, we discuss a setup which features the co-existence of two species of different spins and isospins as the two DM candidates, one being a scalar doublet while the other is a fermion singlet [93]. Thanks to the conversion between the two DM components, we can have a viable parameter region, where one DM component is underabundant and the other component is overabundant or vice versa. However, one cannot obtain the correct relic if both the components are simultaneously underabundant or overabundant. In this model, the low mass region ($m_{\text{DM}} \leq 100$ GeV) is overabundant for both the components. To address this, we add extra DM annihilation channels by adding a charged fermionic singlet under the same $\mathbb{Z}_2$ symmetry as the scalar component. We conclude this work by showing that the entire DM parameter region, i.e., $10 \text{ GeV} \leq m_{\text{DM}} \leq 2 \text{ TeV}$ is compatible with the experimental constraints. The charged fermion, on the other hand, adds extra diagrams in the leptonic final states at LHC, which may substantially change the signal vs background efficiency.

Chapter 3 : The second work in this part [89] discusses another multi-component DM scenario. But its motivation is different from the first [93] in the sense that here we focussed on overcoming the failure of a very popular scalar WIMP DM model (Inert Higgs Doublet Model, or, IDM) with adding a second vector DM particle in the particle spectrum. We show that the DM underabundance over the vast parameter region (100-500 GeV) of IDM can be substantially increased for appropriately chosen interaction strength between the two DM candidates. We also indicate the possible collider probe for the scalar DM component and perform a signal vs background analysis for oppositely signed dilepton + E_T^{miss} final states.

Part II: Coannihilating DM : The standard formalism of WIMP freeze-out is significantly modified if one introduces almost degenerate dark sector particles. Relic density achieved by coannihilation relaxes the parameter space in such scenarios. Here we focus on a second possibility of reviving the shrunk parameter space of WIMPs through *coannihilation* [148] - a phenomenon which is considered an exception in the calculation of thermal abundance in the freeze-out of WIMPs. In the coannihilation phenomena, DM annihilates with another nearly degenerate dark sector particle, and the chemical equilibrium between the annihilating particles ensures the substantial depletion of DM number density. As the direct detection channels are relevant only for $DM \rightarrow q\bar{q}$ interactions, the coannihilation works efficiently to keep the bounds from such measurements at bay. In the two works [90, 91] discussed in this part, we take up the case of quintessential minimal WIMP paradigm of a scalar singlet as the DM candidate. Scalar singlet DM by itself is a very constrained model because of tight direct detection bounds. Here we propose two scenarios, the latter being the logical continuation of the first. Both of these help us in escaping such bound and opening up new possibilities of both rich phenomenologies in relic density analysis and search prospects in the controlled laboratory environment.

Chapter 4 : In the first scenario, [91], we add leptonic coannihilation channels in scalar singlet DM model with an additional vector-like dark lepton doublet. This extra doublet couples with the SM leptonic doublet and the scalar singlet via a non-SM-like Yukawa structure. As a result, (i) since the extra fermionic states interact with other dark sector particles as well as the SM via gauge and Yukawa interactions, it gives rise to new DM annihilation processes including pair annihilation, coannihilation and mediator annihilation channels, and (ii) such a Yukawa structure opens up new production channels for leptonic final states giving much enhancement in cross sections to search for dark matter in the LHC. We showed that a large parameter space from a few tens of GeV up to TeV scale can be allowed by the experimental limits if the scalar DM obtains relic density through coannihilation. Further, depending on the mass splitting between DM and other dark sector particles, we give a quantitative estimate of the transition between mediator annihilation, coannihilation and pair annihilation. We also present the collider probe of this model with multilepton final states and tau-tagging at the LHC using the ROOT based Boosted Decision Tree (BDT) classifier through Toolkit for Multivariate Analysis (TMVA). MVA is necessary for this scenario as the signal background isolation was very difficult using traditional cut-based analysis due to the background almost overshadowing the signal in some cases.

Chapter 5 : In the second scenario, [90], we present a continuation of the previous study [91] with an interesting tweak in the model content. Previously, the gauge invariance between the components of the lepton doublet demanded an introduction of an additional triplet to generate the small mass splitting between DM and the dark sector particles required for coannihilation. Here, in addition to the scalar DM and the lepton doublet, we added an extra vector-like lepton singlet. The mixing between the singlet and a doublet leptonic dark sector produces an arbitrary mass difference between the two components of the doublet in a gauge-invariant way, without introducing any new scale of electroweak symmetry breaking in the theory, making this model more minimalistic. The study provides a useful handle to distinguish between the dark sector particles of different isospins, which is a challenging task otherwise. As the dark leptons coannihilate non-trivially, the mixing effectively enhances the viable parameter space for the relic density constraint. We also show its positive impact in various DM search prospects, e.g., indirect detection and collider searches. In low DM mass regime, our analysis shows that with a non-zero mixing, it is possible to relax the existing indirect search bounds on the upper limit of the DM-Standard Model coupling. In the collider analysis, From the observations of the kinematic distribution of various observables for the $3\tau + E_T^{miss}$ and $\ell\tau + E_T^{miss}$ channels, we conclude that one can clearly distinguish the effects of the mixing parameter that remains unaffected by the change in other free parameters of the model. We also study the statistical significance and present the numerical results signifying the detection prospect of such final states at future LHC luminosities.

Part III: Non-thermal DM : Taking a detour from the thermal WIMP scenario, in this chapter we discuss a completely different theory of DM genesis, known as the freeze-in mechanism [154], which is one of the promising alternatives in non-thermal theories of DM. Here, unlike the WIMP case, DM has very small coupling with SM and its abundance slowly builds up from zero through the decay or scattering of heavier thermal bath particles and saturates to the observed value for proper interaction strength. It's feeble coupling with the visible sector, justifying the name given to this class of particles known as Feebly Interacting Massive Particles (FIMP), keeps it safe from the direct searches. A natural consequence of such a mechanism is the slow decaying partner particles that can leave distinct signatures at the colliders due to their long-lived nature. These particles, depending on their coupling strength, decay within the detector volume leaving displaced decay signatures or can travel beyond the reach of the detectors before decaying, leaving steady charged tracks in the colliders.

Chapter 6 : In the above context, here we have proposed [92] a minimalistic model which can explain DM as both FIMP and WIMP depending on its coupling strength with thermal bath, ie, SM as well as indicates possible long-lived signatures naturally arising from the requirement of the theory. We have systematically identified different areas of the parameter region where DM relic density can be achieved by freeze-in and freeze-out mechanisms. We have considered the minimalistic BSM extension that can satisfy the PLANCK [16] data extending over the thermal and non-thermal regime in the parameter space. This is thanks to the three-body decay which leads to FIMP production from the bath particles for feeble DM-SM interaction limit and on the other hand, coannihilation in the large interaction limit. Based on the range of the parameters involved, our study gives a quantitative estimate of the limiting case where the transition between the thermal and non-thermal regime takes place.

In the context of collider probe, there are several interesting possibilities of long-lived particle (LLP) signatures [3, 7, 163] which naturally arise in the model depending on different mass hierarchies and the coupling strength, and these signatures are discussed in details in the work. We show the possibility of LLP signatures which has not been explored in LHC so far and we indicate the parameter region which can satisfy relic as well as provide these novel exotic signatures. Here it is worth mentioning that the search strategies for LLP are very seriously pursued by particle physics experimentalists recently and dedicated LLP detectors such as FASER [34], MATHUSLA [114] etc. are being proposed or built. Our work, in this context, proposes methods of testing such particles arising naturally in DM scenarios.

Furthermore, we show that the features in the DM phenomenology are generic and can work for a wide class of models of GeV-TeV scale DM and a similar combination of Yukawa couplings representing the DM-SM interaction strength. Lastly, as a natural consequence of the arrangement, this model features right-handed neutrinos that account for the neutrino mass through Type-I see-saw mechanism. This requirement makes the right-handed neutrino long-lived and we indicate the possibilities of testing this in the future LLP detectors.



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