Astro-Particle Aspects of Dark Matter Phenomenology

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To My Family

DECLARATION

I, Mrs. Tanushree Basak, D/o Mr. Barun Kumar Basu, resident of D-7/9, DOS Housing Colony, Vikramnagar, Ahmedabad, 380058, hereby declare that the research work incorporated in the present thesis entitled, "Astro-Particle Aspects of Dark Matter Phenomenology" is my own work and is original. This work (in part or in full) has not been submitted to any University for the award of a Degree or a Diploma. I have properly acknowledged the material collected from secondary sources wherever required. I solely own the responsibility for the originality of the entire content.

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I feel great pleasure in certifying that the thesis entitled, "Astro-Particle Aspects of Dark Matter Phenomenology" embodies a record of the results of investigations carried out by Mrs. Tanushree Basak under my guidance. She has completed the following requirements as per Ph.D regulations of the University.

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I am satisfied with the analysis, interpretation of results and conclusions drawn. I recommend the submission of thesis.

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ABSTRACT

The identity of dark matter (DM) is one of the key outstanding problems in both particle and astrophysics. As the thesis title indicates this work is about identifying a suitable DM candidate and studying its properties. The presence of DM has been supported by a variety of evidence. At galactic and sub-galactic scales, this evidence includes galactic rotation curves, the weak gravitational lensing of distant galaxies by foreground structure, and the weak modulation of strong lensing around individual massive elliptical galaxies. On cosmological scales, observations of the anisotropies in the cosmic microwave background and large scale structure strongly lead us to the conclusion that 80 - 85% of the matter in the universe (by mass) consists of non-luminous and non-baryonic matter. There are several experimental constraints on DM which includes relic density measurement from WMAP and PLANCK, direct detection and indirect detection of DM. The mass and scattering cross-section of the DM off nucleon is probed by direct detection experiments like XENON100, CDMS, DAMA, CoGENT, LUX etc. The indirect detection experiments like PAMELA, AMS02, Fermi-LAT rely on the observations of DM annihilation products such as positrons, antiprotons and photons which might indicate the existence of DM.

In the work presented here we have proposed a supersymmetric extension of Standard Model (SM) with additional hypercharge zero SU(2) triplet and singlet chiral superfields. The triplet sector gives an additional contribution to the scalar masses, and we find that the lightest CP-even Higgs boson can have a mass of 119 - 120 GeV at tree level, and a little radiative correction raises the value to 125 GeV. In this model no significant contributions from stop loops is needed to get the required Higgs mass that alleviates the fine-tuning problem of fixing the stop mass to a high precision at the grand unified theory scale. In the R-parity conserving scenario, we identify the lightest supersymmetric particle (LSP) as the DM candidate. This model naturally gives a neutralino dark matter of mass ~ 100 GeV. We have also explained the 130 GeV γ -ray line, seen at Fermi-LAT, while being consistent with other dark matter observations. We obtain

the required cross-section of 10^{-27} cm³sec⁻¹ for the monochromatic gamma-ray flux through the resonant annihilation of dark matter via pseudoscalar triplet Higgs of mass ~ 260 GeV. The dark matter is predominantly bino-higgsino like, which has large couplings with photons (through higgsino) and gives correct relic density (through bino). In addition, we get the enhanced Higgs diphoton decay rate, $R_{\gamma\gamma} \simeq 1.224$ dominantly contributed by the light chargino-loops, which can account for the reported excess seen in the $h \to \gamma\gamma$ channel by ATLAS.

Another part of this thesis deals with the gauge extension of SM. Here, we adopt the minimal gauged B - L extended SM which contains a singlet scalar and three right-handed neutrinos. The vacuum expectation value of the singlet scalar breaks the $U(1)_{B-L}$ symmetry. Here the third-generation right-handed neutrino is qualified as the dark matter candidate, as an artifact of Z_2 -charge assignment. Relic abundance of the dark matter is consistent with WMAP9 and PLANCK data, only near scalar resonances where dark matter mass is almost half of the scalar boson masses. Requiring correct relic abundance, we restrict the parameter space of the scalar mixing angle and mass of the heavy scalar boson of this model. Besides this, the maximum value of the spin-independent scattering cross section off nucleon is well below the XENON100 and recent LUX exclusion limits and can be probed by future XENON1T experiments. In addition, we compute the annihilation of the dark matter into a two-photon final state in detail and found it to be consistent with the Fermi-LAT upper bound on $\langle \sigma v_{\gamma\gamma} \rangle$ for the Navarro-Frenk-White (NFW) and Einasto profile.

Keywords : Dark Matter, Beyond Standard Model, Supersymmetry Phenomenology, Gauge Extension, Relic Abundance

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List of Abbreviations

DM	Dark Matter
SM	Standard Model
BSM	Beyond Standard Model
WIMP	Weakly Interacting Massive Particles
CDM	Cold Dark Matter
SUSY	Supersymmetry
MSSM	Minimal Supersymmetric Standard Model
NMSSM	Next-to-Minimal Supersymmetric Standard Model
TMSSM	Triplet extended Supersymmetric Standard Model
TSMSSM	Triplet-Singlet extended MSSM
SI	Spin-independent
SD	Spin-dependent
RH	Right-handed
VEV	Vacuum expectation value
EW	Electro-weak
EWSB	Electro-weak symmetry breaking
LHC	Large Hadron Collider
BR	Branching ratio
GC	Galactic Centre

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Chapter 1

Introduction

After the discovery of Standard Model(SM)-like Higgs boson, the last missing particle of the SM, at the LHC [1, 2], a new era of particle physics has started. Although the SM provides very successful and precise description of all experiments in particle physics, it has some theoretical shortcomings. The SM not only suffers from the hierarchy problem, but there remain many open questions like mass of neutrinos; origin of matter-antimatter asymmetry; what is the invisible or 'dark' matter in the universe whose amount is five times the amount of all visible matter etc. Here lies the pressing need to look beyond SM in order to find an explanation of some of these queries. The problem of dark matter is surely one of the most exciting and challenging open questions in physics.

The earliest identification of dark matter came from the velocity dispersions of galaxies within clusters. In 1933, Fritz Zwicky deduced the existence of a non-luminous constituent of the Coma cluster by observing the dynamics of the galaxies contained therein [3], famously conferring upon it the name of "dark matter". The only way the observed velocities of the cluster members could be reconciled with the virial theorem was to postulate that the cluster also contained another large, but unseen, mass component: dark matter. Cosmological observations of the cosmic-microwave-background (CMB) anisotropies, by Wilkinson Microwave Anisotropy Probe (WMAP9) [4], PLANCK satellite [5], constrain the dark-matter density (in units of the critical density) of the Universe to be, $\Omega_{CDM}h^2 = 0.1149 \pm 0.0019$. From the observations we know that, our Universe



Figure 1.1: Rotation curve of dwarf spiral galaxy M33, taken from [6]. The solid curve shows the observed velocity and the dashed curve shows the estimated contribution from the luminous disk.

consists of 71.4% of dark energy, 4.6% of luminous matter and 24% of dark matter. Hence, the non-baryonic dark matter constitutes the majority roughly 80 - 85% of the matter in the Universe. Although, the evidence favouring the existence of DM is extremely compelling, its nature mostly remains unknown. In this thesis, we focus on the extensions of SM and propose various particle physics models to explain experimental consequences of DM. In the pursuit of understanding DM phenomenology many ideas from particle physics have been used, and in equal measure experimental observations from astrophysics and cosmology have been used to constrain ideas in particle physics. This interplay between theoretical particle physics and astrophysical observations lies at the core of the work presented here.

1.1 The evidence for dark matter

Rotation Curves : The earliest and as yet the most convincing evidence in support of the existence of DM came from the rotation curves [7] of galaxies (graph of circular velocities of luminous gas in a galaxy as a function of their distance from the galactic centre). According to Newtonian dynamics the radial velocity of a galaxy in a cluster is given by,

$$v(r) = \sqrt{\frac{GM(r)}{r}} \tag{1.1}$$

where, $M(r) = \int 4\pi \rho(r) r^2 dr$ and $\rho(r)$ is the density profile. The above equation shows that velocity should be falling as $1/\sqrt{r}$, but instead the rotation curve in Fig.1.1 shows a considerably flatter behaviour, which suggests the existence of an halo with $M(r) \propto r$. Therefore, these rotation curves pointed out that the visible mass in those galaxies could not account for the observed circular velocities and hence postulated the presence of a large unseen component of mass inside the galaxy.

Gravitational lensing : According to General Relativity, the presence of mass causes the space in its vicinity to curve. Clusters, galaxies and stars are massive enough to bend and focus light rays passing through their gravitational potential, which acts as a lens. As a result light from distant galaxies, quasars and stars are gravitationally lensed [8] by other clusters and galaxies which lie in their path. The amount of lensing depends on the mass of the object (the lens) causing this effect. It can therefore be used to determine the mass of astrophysical objects ranging from planets to galaxies and galaxy clusters.

A different example of gravitational lensing in the context of DM is that of the so called Bullet cluster [10], shown in Fig.1.2. This figure shows the collision between a smaller cluster (bullet) with a primary cluster. The distribution of mass from the weak gravitational lensing (blue region) and that from the X-ray map (pink region), which consists of mostly baryonic matter, shows that the majority of matter in the two galaxies is non-baryonic and it is primarily in halo region of the galaxy. Moreover, the separation in the two regions shows that while the gas clouds in the two galaxies exerted friction on each other resulting in the bullet shape of the right most cluster, the DM halos of the two galaxies pass through each other and the gas clouds without undergoing any collision. This strongly indicates the presence of a collisonless, non-baryonic DM halo.

Cosmological evidence : The universe displays a very complex structure on a large scale. Galaxies of stars are part of a cluster, clusters of such galaxies



Figure 1.2: Composite image of the matter inside the galaxy cluster 1E 0657-56, also known as the "bullet cluster". The blue region shows the lens- ing map while the pink region shows the X-ray data associated with the gas clouds. The clear separation between the X-ray and the lensing maps shows that most of the matter inside the galaxies is collisionless dark matter. X-ray : NASA/CXC/CfA/Magellan/U.Arizona [9], Optical data : NASA/STScI [10] and Lensing maps : NASA/STScI; ESO/WFI/Magellan/U.Arizona [10].

are again part of superclusters which are then arranged into large scale sheets, filaments and voids. Presumably, the pattern of galactic superstructure reflects the history of gravitational clustering of matter since the Big Bang. If dark matter were present during structure formation, it should have influenced the pattern of large-scale structure we see today. Large scale structure surveys like the 2DFGRS [11] and the SDSS [12] can also provide information on the total matter density in the universe [13, 14]. Large-scale cosmological 'N-body' simulations [15–18] demonstrate that the observed large-scale structure of luminous matter could only have been formed in the presence of a substantial amount of dark matter. A recent analysis in [14] indicates a total matter density of $\Omega_M = 0.29$. The Big bang nucleosynthesis (BBN) data gives the baryonic density to be $\Omega_B = 0.04$. Combining this with the BBN result gives the non-baryonic DM density to be $\Omega_{CDM} = 0.25$.

1.2 Constraints on Dark matter candidate

A good DM particle candidate should fulfill a series of important properties [19–22] in order to provide a convincing explanation to all the observed phenomenology:

- It should be weakly interacting to ordinary matter, i.e, SM particles and electrically neutral, i.e. with neither electromagnetic nor strong interactions.
- It should be long-lived with lifetime larger than H_0^{-1} so that it survives since the time in the early Universe, when they were created.
- It has to be cold, i.e, it should be non-relativistic when it decouples from the radiation in order to not erase the density fluctuations at galaxy scales. At most, the dark matter may be warm, with free-streaming lengths on the order of cluster scales of a few Mpc.
- It must be massive enough to account for the measured Ω_{DM} .
- It must be consistent with observations (BBN, relic density) and present constraints (direct and indirect detections).

1.3 Dark Matter particle candidates

The concept of dark matter does not find an explanation in the framework of the Standard Model (SM). Plenty of extensions of the SM were put forward with a motivation to introduce a suitable DM candidate. Potentially the only indication compatible with cosmological measurements is that dark matter is composed of non-baryonic, neutral and weakly interacting particles. Weakly interacting massive particles (WIMPs) are the most favored and widely-studied DM candidate, as they satisfy the astrophysical and cosmological criteria, and offer the possibility of detectable experimental signals. Examples of WIMPs include the lightest neutralino in supersymmetry (see Sec.1.5.1), the lightest Kaluza-Klein (KK) particle [23, 24], a scalar dark matter (see Sec.1.5.2), an additional inert Higgs boson [25–28], RH-neutrino dark matter (see Chap.4) etc.

In the literature, several non-WIMPs candidates (for review see [21, 22, 29]) were also proposed, among which the most relevant are: Sterile neutrinos, Axions, Gravitino, Axino, Dark matter from little Higgs model, Superheavy dark matters or Wimpzillas.

1.4 The WIMP Miracle

The most intriguing piece of cosmological evidence in favor of WIMP dark matter is that thermally produced WIMPs naturally have a relic abundance close to that observed for dark matter. In the early universe when the temperature was high enough $(T \gg m_{\chi})$, the DM particles were in thermal equilibrium with the rest of the cosmic plasma. In order for this particle to remain in thermal equilibrium, it should interact sufficiently with its surrounding. The evolution of the WIMP number density n_{χ} is given by the Boltzmann equation,

$$\frac{dn_{\chi}}{dt} + 3Hn_{\chi} = \langle \sigma v \rangle (n_{\chi,eq}^2 - n_{\chi}^2)$$
(1.2)

where H is the Hubble expansion rate, $\langle \sigma v \rangle$ is the thermally averaged total annihilation cross-section, and $n_{\chi,eq}$ is the equilibrium WIMP number density. As the universe expands and its temperature falls, the number density of WIMPs decreases exponentially. Thus, the annihilation rate reduces and becomes smaller than the Hubble expansion rate. Then the DM species decouples from the cosmic plasma and number density experiences a "freeze-out" - hence we observe a significant relic abundance of DM today. When freeze-out occurs, $\Gamma(T_f) \simeq$ $n(T_f)\langle \sigma v \rangle(T_f) \sim H(T_f)$. Thus, WIMPs freeze out when they are nonrelativistic, at a temperature where their equilibrium number density is Boltzmann suppressed and their velocity is small.

The WIMP relic abundance is then given by the present-day WIMP density [30],

$$\Omega_{\chi}h^2 = 1.1 \times 10^9 \frac{x_f}{\sqrt{g^*}m_{Pl}\langle\sigma v\rangle_{ann}} \text{GeV}^{-1} , \qquad (1.3)$$

where $x_f = m_{\chi}/T_f$ with T_f as freeze-out temperature. m_{Pl} is Planck mass = 1.22×10^{19} GeV, and, g^* is effective number of relativistic degrees of freedom (we use, $g^* = 100$ and $x_f = 20$). It can be expressed in a more common form as,

$$\Omega_{\chi} h^2 \approx 0.1 \left(\frac{3 \times 10^{-26} \text{cm}^3/\text{sec}}{\langle \sigma v \rangle_{ann}} \right)$$
(1.4)

Now, if we calculate the annihilation cross-section of a WIMP, having mass in the weak scale range $(m_{\chi} \sim M_W)$, resulting from a weak-interaction, we obtain, $\sigma \sim \alpha^2/m_{\chi}^2$. Here, α is the fine-structure constant. Thus, we get, $\sigma \approx 2 \text{pb} \left(\frac{100 \text{GeV}}{m_{\chi}}\right)^2$. Hence, a weak-scale annihilation cross section naturally gives a thermally produced WIMP relic abundance that matches the observed dark-matter relic abundance -this striking coincidence is known as the "WIMP miracle".

1.5 Models for dark matter

1.5.1 Supersymmetric Dark Matter

The most popular and exhaustive extension of SM is Supersymmetry (SUSY) (for reviews see [31–34]), which overcomes not only many shortcomings of the SM but has also additional and very attracting features. This is a symmetry relating fermions to bosons such that for each fermionic degree of freedom there is a bosonic degree of freedom. This extends the particle content of the Standard Model (SM) such that each particle in the SM has a corresponding superpartner (or partners). SUSY solves the hierarchy problem in a very simple and elegant way: since every divergent loop diagram containing a SM fermion is matched by corresponding scalar sfermion loop diagrams.

The particle content in the Minimal supersymmetric Standard Model (MSSM) [35–41] is the same as of the SM plus the superpartners and two Higgs doublets (instead of one as in the SM). Two Higgs doublets are needed for anomaly cancellation and to give mass to both up- and down-type quarks and will result in five physical Higgs bosons. If supersymmetry were unbroken, a SM particle and its superpartner would have the same mass and quantum numbers (except for spin). Since we haven't seen these particles, we can conclude that supersymmetry is broken at the energies probed by present accelerators. The superpotential of the MSSM is given by,

$$\mathcal{W} = \mu H_u \cdot H_d + y_u Q_L \cdot H_u U_R + y_d Q_L \cdot H_d D_R + y_e L_L \cdot H_d E_R \tag{1.5}$$

The scalar potential suggests that the bound on lightest Higgs boson mass at tree level is, $m_h \leq M_z |\cos 2\beta|$, which has been exceeded by LEP and LHC, where $m_h \geq 114.4$ GeV and 125 - 126 GeV respectively. Therefore, a significant loop correction with maximal top-stop mixing is required to raise m_h upto 126 GeV.

Also an additional discrete symmetry called R-parity is defined in the MSSM to evade proton-decay as, $R = (1)^{2S+L+3B}$, with S, L and B respectively the spin, leptonic and baryonic quantum number. We see that all SM particles have a parity of +1 while all supersymmetric particles have a parity of -1. If the R-parity is conserved throughout the interactions - i.e. if we impose this conservation by forcing a symmetry on the lagrangian, then the Lightest Supersymmetric Particle (LSP) - will be stable and will not decay into lighter normal matter. This is a good way to ensure the stability at cosmological scales of dark matter.

The neutralinos are linear combinations of the superpartners of the neutral gauge bosons (i.e, gauginos) and the Higgs bosons (i.e, higgsinos). In the gauge-basis $(\tilde{B}, \tilde{W^0}, \tilde{H^0_d}, \tilde{H^0_u})$, the mass matrix is given by,

$$\mathcal{M}_{\bar{G}} = \begin{pmatrix} M_1 & 0 & -\cos\beta\sin\theta_w M_Z & \sin\beta\sin\theta_w M_Z \\ 0 & M_2 & \cos\beta\cos\theta_w M_Z & -\sin\beta\cos\theta_w M_Z \\ -\cos\beta\sin\theta_w M_Z & \cos\beta\cos\theta_w M_Z & 0 & -\mu \\ \sin\beta\sin\theta_w M_Z & -\sin\beta\cos\theta_w M_Z & -\mu & 0 \end{pmatrix}$$

where, M_1 , M_2 are the soft breaking mass parameters for Bino and Wino respectively. Therefore, the neutralino mass matrix can be diagonalized analytically to give the four neutralinos,

$$\tilde{\chi}_{i}^{0} = N_{i1}\tilde{B} + N_{i2}\tilde{W}_{3}^{0} + N_{i3}\tilde{H}_{d}^{0} + N_{i4}\tilde{H}_{u}^{0}, \qquad (1.6)$$

the lightest of which, $\tilde{\chi}_1^0$ serves as the main candidate for dark matter in SUSY models. Despite of all the success, MSSM suffers from some intrinsic problems.

Going beyond MSSM, provides a solution to the μ -problem, by simply adding a singlet superfield, i.e the so-called NMSSM models. Here also the neutralino is still the DM, but its properties can be quite different due to the contributions of a gauge singlet. Other possible extension of MSSM are by either adding new particles like triplets or extending the gauge group. All of these extensions of the MSSM have significant influence on the properties of the dark matter particle. There are also other well-motivated SUSY dark matter candidates like gravitino, axino, sneutrino etc.

1.5.2 Gauge Singlet Scalar Dark Matter

The scalar singlet extension of SM [42–47] is the most simplified Higgs-portal model to account for a WIMP candidate. The real singlet S', stabilized by odd \mathbb{Z}_2 -parity, acts as a viable DM candidate. It interacts only with the SM Higgs boson through the renormalizable interaction term present in the lagrangian,

$$\mathcal{L} = \mathcal{L}_{SM} + \frac{1}{2} (\partial S')^2 - \frac{1}{2} \mu_{S'}^2 S'^2 + \mathcal{L}_{int} - \lambda S'^4$$
(1.7)

where, $\mathcal{L}_{int} = -\lambda_{S'} |\Phi|^2 S'^2$ denotes the interaction between the SM-Higgs and dark matter.

The mass of the DM after EWSB becomes, $m_{DM}^2 = \mu_{S'}^2 + \frac{1}{2}\lambda_{S'}v^2$. The coupling between DM and SM-Higgs, i.e, $\lambda_{S'}$ is constrained from the invisible decay width of Higgs boson when $m_{DM} \leq m_h/2$, such that BR $(h \rightarrow SS) \leq 0.20$. The dark matter annihilates through SM-Higgs into SM-particles and thus account for the correct relic abundance, when $m_{DM} \sim m_h/2$. The coupling and scattering cross-section can also be constrained from the direct detection experiments.

1.5.3 Singlet Fermionic Dark Matter

The singlet fermionic dark matter (SFDM) model is a renormalizable extension of SM with a hidden sector containing a scalar singlet Φ_s and a singlet Dirac fermion ψ [48–53]. Here, the singlet fermionic dark matter ψ , interacts with the SM sector via the singlet Φ_s which mixes with the SM-Higgs doublet Φ . Therefore, this is also an example of Higgs-portal model. The lagrangian of the SFDM model is given as,

$$\mathcal{L} = \mathcal{L}_{SM} + \mathcal{L}_{hid} + \mathcal{L}_{int} \tag{1.8}$$

where,

$$\mathcal{L}_{hid} = \mathcal{L}_{\Phi_s} + \bar{\psi}(i\partial_\mu\gamma^\mu - m_\psi)\psi - \lambda_{\psi S}\,\bar{\psi}\psi\Phi_s \tag{1.9}$$

$$\mathcal{L}_{int} = \frac{\lambda_1'}{2} \Phi^{\dagger} \Phi \Phi_s + \frac{\lambda_2'}{2} \Phi^{\dagger} \Phi \Phi_s^2$$
(1.10)

and

$$\mathcal{L}_{\Phi_s} = \frac{1}{2} (\partial \Phi_s)^2 - \frac{m_{\Phi_s}^2}{2} \Phi_s^2 - \frac{\lambda'}{3} \Phi_s^3 - \frac{\lambda''}{4} \Phi_s^4$$
(1.11)

After EWSB, the singlet field Φ_s can be written as, $\Phi_s = x + s$, where x is the VEV of Φ_s and SM-Higgs doublet Φ is same as in eq.4.3. The two scalar eigenstates are denoted as,

$$h_2 = \sin\theta \, s + \cos\theta \, \phi \tag{1.12}$$

$$h_1 = \sin\theta \phi - \cos\theta s \tag{1.13}$$

where, h_2 is identified as the SM-Higgs boson and $\cos \theta(\sin \theta)$ is the scalar-mixing. Now, the mass of the DM is given by, $m_{DM} = m_{\psi} + \lambda_{\psi S} x$, with m_{ψ} as a free parameter. The DM interaction strength depends on the parameter $\lambda_{DM} = \lambda_{\psi S}$. Here, the scalar mixing angle and DM-coupling are subject to various constraints like LHC bound on SM-Higgs boson, relic abundance of DM and upper bound on the DM-nucleon scattering cross section.

1.6 Dark matter Searches

1.6.1 Direct Detection

In direct experiments one looks for these WIMPs passing through a detector and scattering off some nucleus. A variety of detectors [54–64] designed to be sensitive to the nuclear recoils induced by collisions with WIMPs are currently



Figure 1.3: Spin-independent elastic WIMP-nucleon cross section (σ_p^{SI}) as a function of WIMP mass (m_{χ}) (taken from [54]). The thick blue line shows the 90% CL for the latest XENON100 data. The limits from CDMS(dashed orange line), XENON(2010) (thin black line) and EDELWEISS (dotted pink line) are also shown. And finally the 90% CL regions for CoGeNT (green) and DAMA (red, without channeling) are shown.

collecting data, and have placed bounds on the WIMP-nucleon cross section – WIMP-mass parameter space. This scattering can be detected and, if found, would be an evidence for WIMPs in the galactic halo. These studies also depend on astrophysical input like in particular, the local phase-space distribution of dark-matter particles.

Exclusion limits placed upon the WIMP mass and spin-independent nuclear scattering cross-sections by various direct detection experiments are shown in Fig.1.3. Among these experiments, however, XENON 100 sets the most stringent upper limit on the spin-independent (SI) scattering cross-section of DM off the nuclei, which is of the order of ~ 2 × 10⁻⁴⁵ cm² for $m_{\chi} \simeq 55$ GeV. Recently, the LUX experiment sets a minimum upper bound on the SI- cross section of 7.6 × 10⁻⁴⁶ cm² (shown in Fig. 1.4) at a WIMP mass of 33 GeV – this result is



Figure 1.4: The LUX 90% confidence limit on the spin-independent elastic WIMP-nucleon cross section (blue), together with the $\pm 1\sigma$ variation from repeated trials (taken from [57]), where trials fluctuating below the expected number of events for zero BG are forced to 2.3 (blue shaded). We also show Edelweiss II [64] (dark yellow line), CDMS II [62] (green line), XENON100 100 live-day [55] (orange line), and 225 live-day [56] (red line) results. The inset (same axis units) also shows the regions measured from annual modulation in CoGeNT [60] (light red, shaded) and DAMA/LIBRA allowed region [58].

in strong disagreement with the earlier experimental bounds.

1.6.2 Indirect Detection

In the indirect searches, one looks not for the WIMPs directly, but for signals coming from annihilation of two WIMPs. For example, their annihilations in the halo will result in a γ - and anti-proton-flux which can be searched for. A wide array of cosmic-ray and gamma-ray observatories – both in space and on the ground – are currently searching for indirect signals. There are many hints of the DM annihilation in the high energy cosmic ray spectrum of positrons and γ -



Figure 1.5: The positron fraction compared with the most recent measurements from AMS-02 [67], PAMELA [65, 66] and Fermi-LAT [68].

rays. DM annihilation may account for the excess positron flux seen in PAMELA [65, 66] and AMS-02 [67] experiments, shown in Fig. 1.5. Observation of γ -rays in Fermi-LAT [68], Hess [69] give possible signals of DM annihilation into γ -rays.

Whether searching for annihilations or decays, the most promising targets are those with large dark matter densities and/or low astrophysical backgrounds. The Galactic Center (GC) would seem the most obvious target given its distance and dark matter concentration, but it is also one of the most difficult areas to work with because of its complex and poorly-understood background. Two years back, a tantalizing hint of DM was found in the analysis of the Fermi-LAT gamma-ray data [68] which revealed the existence of a peak at around 130 GeV coming from the vicinity of the GC.

1.6.3 Collider search on Dark Matter

Another possibility of searching for dark matter is through the production of DM in collider experiments from the annihilation of SM-particles at sufficiently high energies. Any WIMPs produced at colliders will escape from the detector



Figure 1.6: Observed 90% C.L. upper limits on the χ -nucleon scattering cross section as a function of m_{χ} for spin-independent effective operators mediating the interaction of the dark-matter particles with the $q\bar{q}$ initial state (taken from [70]). The limits are compared with results from the published ATLAS hadronically decaying W/Z and $j + \chi \chi$ searches, CoGeNT [59], XENON100 [56], CDMS [61, 63], and LUX [57].

unnoticed. The most obvious collider WIMP signature is expected to be missing transverse energy (missing ET), which refers to an apparent missing component of the total final-state momentum in the direction transverse to a collider beam. Monojet searches typically give the strongest constraints [71]. However, one can hope to see DM also in other channels as, monophoton searches, mono-W searches and mono-Z searches.

In the frame-work of Higgs portal dark matter, bounds from invisible Higgs decays can be translated into bounds on the WIMP scattering cross-section[76]. The resulting bounds apply only for $m_{\chi} < m_h/2$ and depend on whether the WIMP is a scalar, fermion or vector [77]. For effective operators inducing spinindependent interactions, LHC searches [70] are typically inferior to direct de-


Figure 1.7: Observed 90% C.L. upper limits on the χ -nucleon scattering cross section as a function of m_{χ} for the spin-dependent D9 effective operators mediating the interaction of the dark-matter particles with the $q\bar{q}$ initial state (taken from [70]). The limits are compared with results from the published ATLAS hadronically decaying W/Z and $j + \chi \chi$ searches, COUPP [72], SIMPLE [73], PICASSO [74], and IceCube [75].

tection (except for very light masses), since the latter benefit from coherent enhancement (shown in Fig.1.6). But for spin-dependent interactions, direct detection cross sections are not enhanced and LHC searches typically give the strongest bounds, as shown in Fig.1.7.

1.7 Thesis overview

The thesis is organized as follows – Chap.2 is about the newly proposed tripletsinglet extension of Minimal Supersymmetric Standard Model (TSMSSM), where we have extended the superpotential of MSSM with one singlet and hypercharge zero SU(2) triplet chiral superfields. The contributions from the triplet and singlet helps in enhancing the lightest physical Higgs mass at tree-level and thus reduces the fine-tuning compared to other SUSY models. This model naturally accommodates a neutralino-LSP dark matter candidate in the mass range \sim 100 GeV. We will discuss in detail the phenomenological aspects of this model including dark matter, one-loop correction to Higgs mass and improvement of fine-tuning.

In Chap.3 we will address the recent observations of monochromatic gammaray line-like feature in the context of TSMSSM. We perform a scan over the parameter space of the TSMSSM model, and choose a suitable benchmark set which satisfies all phenomenological requirements. We will show that the required cross-section for explaining the gamma-ray spectral feature can be achieved via resonant enhancement in the dark matter annihilation cross-section into two photon. Futher, we will show that the reported excess in the $h \to \gamma \gamma$ decay rate can be also explained.

Next in Chap.4, we will focus on the minimal $U(1)_{B-L}$ extension of Standard model which contains a right-handed (RH) neutrino dark matter candidate as an artifact of Z_2 charge assignment. We will also show that the mass of the heavy scalar boson and the scalar mixing angle can be constrained in order to obtain correct relic abundance of dark matter. We further discuss the annihilation of such DM into two-photon final state.

We will summarize and present the impact of our work in the last chapter. We will also point out some of the recent observations like (i)3.55 keV X-ray line emission, (ii) a spatially extended excess of $\sim 1-3$ GeV gamma rays from the regions surrounding the galactic center, (iii) positron excess seen by AMS-02, PAMELA etc – which are likely to be addressed in the context of DM in future.

Chapter 2

Dark matter in the Triplet-singlet extension of MSSM

2.1 Overview

The ATLAS and CMS collaborations [1, 2] have narrowed down the allowed range of the SM-like Higgs mass to the region 124-126 GeV. A light Higgs is favored in supersymmetry although the MSSM predicts a tree level upper bound on the lightest CP-even Higgs mass as $m_h < M_Z \cos 2\beta$. Within MSSM, loop corrections can give required large corrections to Higgs mass provided the stop is heavier than 1 TeV or there is near maximal stop mixing. Implications of the 125 GeV Higgs for the MSSM and constrained-MSSM parameter space have been extensively studied [35–41, 78–82]. Going beyond MSSM, in order to get a larger tree-level Higgs mass, the simplest extension is a singlet superfield in the NMSSM model [83–91]. The singlet interaction with the two Higgs doublet of MSSM is via the $\lambda SH_u \cdot H_d$ term. The Higgs mass is now given by the relation $m_h^2 = M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta + \delta m_h^2$, where δm_h^2 is due to radiative-correction. Taking $\lambda = 0.7$ (larger values would make it flow to the non-perturbative regime much below the GUT scale) and $\tan \beta = 2$ the radiative correction needed to get a 125 GeV Higgs mass is $\delta m_h = 55$ GeV which is an improvement over the $\delta m_h = 85 \text{GeV}$ needed in the MSSM. However fine-tuning of the stop mass is still

required in NMSSM to get the required Higgs mass [87].

Also by extending the MSSM gauge group in a suitable way, the new Higgs sector dynamics can push the tree-level mass well above the tree-level MSSM limit if it couples to the new gauge sector [92–97]. In most of the cases the nondecoupling D-terms contribute non-trivially to increase the tree-level mass of the SM-like Higgs boson. Recent analysis of the SUSY model based on $SU(3)_C \times$ $SU(2)_L \times U(1)_R \times U(1)_{B-L}$ gauge group [98] has shown that the tree level physical Higgs mass can be atmost 110 GeV and through the one-loop correction it can be raised considerably. Another recent work on MSSM extended by a U(1) gauged Peccei-Quinn symmetry [97] where the new D-terms can raise the tree-level mass well enough to accommodate the 125 GeV Higgs boson without significant radiative correction and hence requires less fine-tuning.

An important aspect of the 125 GeV Higgs mass is that the parameter space of thermal relic for dark matter is severely restricted. In MSSM, the LSP is a Higgsino at the TeV scale [81]. In NMSSM, SUSY partner of the singlet scalar -the singlino mixes with the neutralinos to provide a light dark matter [38, 99]. To our knowledge, the dark matter in triplet-extended MSSM has not been studied so far. The extension of MSSM by extending it with a Y = 0and $Y = 0, \pm 1$ SU(2) triplet superfields has been studied [100–102] where the tree level contribution to the Higgs mass from the triplet Higgs sector has been estimated. It has been shown in [102] that with the Y = 0 triplet superfield the tree-level Higgs mass can be raised to 113 GeV which would still require substantial loop corrections from stops. Recently, the MSSM extended by two real triplets $(Y = \pm 1)$ and one singlet [103] has been studied with a motivation to solve the μ -problem as well as to obtain a large correction to the lightest Higgs mass. The analysis of the dark matter sector of this model will be complicated as the LSP will be the lightest eigenstate of the 7×7 neutralino mass matrix which has not yet been done. In this chapter, we explore the minimal extension of the MSSM which can give a tree-level Higgs mass of 119-120 GeV. We find that by extending the MSSM by adding a singlet and a Y = 0 SU(2) triplet superfields, this aim can be achieved. The upper bound on the tree-level mass of the lightest CP-even Higgs is given in equation (2.27). With this tree level Higgs mass the stop mass need not be very heavy and this solves the fine-tuning problem of the Higgs mass in MSSM and NMSSM [87]. We also study the dark matter candidates in this model which is obtained by diagonalizing the 6 × 6 neutralino mass matrix. We find a viable dark matter with mass 100 GeV, which is a mixture of the Higgsino and Triplino (the fermionic partner of the neutral component of the triplet Higgs). We fix two sets of benchmark parameters at the electroweak scale which would give a the 125 GeV and dark matter relic density $\Omega h^2 = 0.1109 \pm 0.0056$ compatible with WMAP-7 measurements [104]. We find that the direct detection cross section of the dark matter is $\sigma_{SI} \simeq 10^{-43} cm^2$, which is compatible with the direct detection experiments like XENON100 [105].

This chapter is organized as follows : In the next section, we discuss the superpotential of our model and we derive the scalar potential from the D-terms and F-terms and from the various soft-breaking terms. In section (2.3) we give a detailed analysis of the Higgs sector and we calculate the CP-even, CP-odd and Charged Higgs mass matrices. In section (2.4), the neutralino and the chargino mass matrices are discussed. The numerical results based on this model are discussed in detail in section (2.5). We show the results for two sets of benchmark points which include the parameters like couplings, tri-linear soft breaking terms, soft masses and the fermionic and scalar mass spectrum. In section (2.5.2) we discuss the Dark Matter from the neutralino sector of this model and its phenomenology, which is one of the main results of this chapter. We have also taken into account the one-loop corrections to the lightest physical Higgs mass (see section 2.6) and shown a quantitative improvement of the level of fine-tuning compared to other models in section 2.7.

2.2 Model

In this model [106], by taking naturalness of the Higgs mass as a guiding criterion, we have extended the superpotential of the minimal supersymmetric standard model by adding one singlet chiral superfield S and one SU(2) triplet chiral superfields T_0 with hypercharge Y = 0. The most general form of the superpotential for this singlet-triplet extended model can be written as,

$$\mathcal{W} = (\mu + \lambda \hat{S}) \hat{H}_{d} \cdot \hat{H}_{u} + \frac{\lambda_{1}}{3} \hat{S}^{3} + \lambda_{2} \hat{H}_{d} \cdot \hat{T}_{0} \hat{H}_{u} + \lambda_{3} \hat{S}^{2} Tr(\hat{T}_{0}) + \lambda_{4} \hat{S} Tr(\hat{T}_{0} \hat{T}_{0}) + W_{Yuk.} , \qquad (2.1)$$

where, $\hat{H}_{u,d}$ are the Higgs doublets of the MSSM and the Yukawa superpotential W_{Yuk} is given as,

$$W_{Yuk.} = y_u \hat{Q}_L \hat{H}_u \hat{U}_R + y_d \hat{Q}_L \hat{H}_d \hat{D}_R + y_e \hat{L}_L \hat{H}_d \hat{E}_R$$
(2.2)

In terms of the components, we have

$$\hat{H}_u = \begin{pmatrix} \hat{H}_u^+ \\ \hat{H}_u^0 \end{pmatrix}, \ \hat{H}_d = \begin{pmatrix} \hat{H}_d^0 \\ \hat{H}_d^- \end{pmatrix} \text{ and } \hat{T}_0 = \begin{pmatrix} \frac{\hat{T}^0}{\sqrt{2}} & -\hat{T}_0^+ \\ \hat{T}_0^- & \frac{-\hat{T}^0}{\sqrt{2}} \end{pmatrix}$$

Here, $(\hat{T}_0^-)^* \neq -\hat{T}_0^+$, which would not have been true for real Higgs triplet in non-supersymmetric models. We can solve the μ -problem by starting with a scale invariant superpotential, given as

$$W_{sc.inv.} = \lambda \hat{S} \hat{H}_d \hat{H}_u + \frac{\lambda_1}{3} \hat{S}^3 + \lambda_2 \hat{H}_d \hat{T}_0 \hat{H}_u + \lambda_4 \hat{S} Tr(\hat{T}_0 \hat{T}_0) + W_{Yuk.}$$
(2.3)

where the SU(2) invariant dot product is defined as,

$$\hat{H}_{d} \hat{T}_{0} \hat{H}_{u} = \frac{1}{\sqrt{2}} (\hat{H}_{d}^{0} \hat{T}^{0} \hat{H}_{u}^{0} + \hat{H}_{d}^{-} \hat{T}^{0} \hat{H}_{u}^{+}) - (\hat{H}_{d}^{0} \hat{T}_{0}^{-} \hat{H}_{u}^{+} + \hat{H}_{d}^{-} \hat{T}_{0}^{+} \hat{H}_{u}^{0}) \quad (2.4)$$

This superpotential (2.3) also has an accidental Z_3 -symmetry, i.e. invariance of the superpotential on multiplication of the chiral superfields by the factor of $\frac{2\pi i}{3}$. By, this choice we are eliminating the μ -parameter but an effective μ -term is generated when the neutral components of S and T_0 acquire vev's v_s and v_t respectively,

$$\mu_{eff} = \lambda v_s - \frac{\lambda_2}{\sqrt{2}} v_t \tag{2.5}$$

Therefore, in terms of the neutral components of the super-fields the equation (2.3) sans W_{Yuk} can be re-written as,

$$W^{neu} = -\lambda \hat{S} \hat{H}_u^0 \hat{H}_d^0 + \frac{\lambda_1}{3} \hat{S}^3 + \frac{\lambda_2}{\sqrt{2}} \hat{H}_d^0 \hat{T}^0 \hat{H}_u^0 + \lambda_4 \hat{S} \hat{T}^0 \hat{T}^0$$
(2.6)

2.2.1 Scalar potential

The scalar potential involving only Higgs field can be written as,

$$V = V_{SB} + V_F + V_D \tag{2.7}$$

In the above equation, V_{SB} consists of the soft-supersymmetry breaking term associated with the superpotential in equation(2.3), is given by

$$V_{SB} = m_{H_u}^2 [|H_u^0|^2 + |H_u^+|^2] + m_{H_d}^2 [|H_d^0|^2 + |H_d^-|^2] + m_S^2 |S|^2 + m_T^2 Tr(T_0^\dagger T_0) + (-\lambda A_\lambda SH_u H_d + \frac{\lambda_1}{3} A_{\lambda_1} S^3 + \lambda_2 A_{\lambda_2} H_d T_0 H_u + \lambda_4 B_\lambda STr(T_0^2) + h.c)$$
(2.8)

In equation(2.7) V_F is the supersymmetric potential from F-terms, given by

$$V_{F} = |-\lambda S H_{d}^{0} + \frac{\lambda_{2}}{\sqrt{2}} H_{d}^{0} T^{0} - \lambda_{2} H_{d}^{-} T_{0}^{+}|^{2} + |-\lambda S H_{u}^{0} + \frac{\lambda_{2}}{\sqrt{2}} H_{u}^{0} T^{0} - \lambda_{2} H_{u}^{+} T_{0}^{-}|^{2} + |\frac{\lambda_{2}}{\sqrt{2}} (H_{u}^{0} H_{d}^{0} + H_{d}^{-} H_{u}^{+}) + 2\lambda_{4} S T^{0}|^{2} + |\lambda (H_{d}^{-} H_{u}^{+} H_{u}^{0} H_{d}^{0}) + \lambda_{1} S^{2} + \lambda_{4} (T^{0^{2}} - 2T_{0}^{+} T_{0}^{-})|^{2} + |\lambda S H_{d}^{-} + \frac{\lambda_{2}}{\sqrt{2}} T^{0} H_{d}^{-} - \lambda_{2} H_{d}^{0} T_{0}^{-}|^{2} + |-\lambda_{2} H_{d}^{-} H_{u}^{0} - 2\lambda_{4} S T_{0}^{-}|^{2} + |\lambda S H_{u}^{+} + \frac{\lambda_{2}}{\sqrt{2}} T^{0} H_{u}^{+} - \lambda_{2} H_{u}^{0} T_{0}^{+}|^{2} + |-\lambda_{2} H_{u}^{+} H_{d}^{0} - 2\lambda_{4} S T_{0}^{+}|^{2}$$
(2.9)

whereas the F-term for the neutral scalar potential can be derived from equation (2.6) as,

$$V_{F_{neu}} = \sum_{i} \left| \frac{\partial W_{scalar}^{neu}}{\partial \phi_i^0} \right|^2$$
(2.10)

where, ϕ_i^0 stands for H_u^0, H_d^0, S, T^0 and W_{scalar}^{neu} is the scalar counter-part of the neutral superpotential W^{neu} .

Finally, V_D is supersymmetric potential from D-terms in equation (2.7), given by

$$V_{D} = \frac{g_{1}^{2}}{8} [|H_{d}^{-}|^{2} + |H_{d}^{0}|^{2} - |H_{u}^{+}|^{2} - |H_{u}^{0}|^{2}]^{2} + \frac{g_{2}^{2}}{8} [|H_{d}^{-}|^{2} + |H_{d}^{0}|^{2} - |H_{u}^{+}|^{2} - |H_{u}^{0}|^{2} + 2|T_{0}^{+}|^{2} - 2|T_{0}^{-}|^{2}]^{2} + \frac{g_{2}^{2}}{8} [H_{d}^{0*}H_{d}^{-} + H_{u}^{+*}H_{u}^{0} + \sqrt{2}(T_{0}^{+} + T_{0}^{-})T_{0}^{*} + h.c]^{2} - \frac{g_{2}^{2}}{8} [H_{d}^{-*}H_{d}^{0} + H_{u}^{0*}H_{u}^{+} + \sqrt{2}(T_{0}^{+} - T_{0}^{-})T_{0}^{*} + h.c]^{2}$$
(2.11)

2.2.1.1 EWSB

After Electroweak symmetry breaking, only the neutral components of the scalars fields acquire vev's, i.e,

$$\langle H_u^0 \rangle = v_u , \langle H_d^0 \rangle = v_d , \langle S \rangle = v_s \text{ and } \langle T^0 \rangle = v_t$$

The neutral-scalar part of the chiral superfields can be decomposed into real and imaginary parts,

$$H_u^0 = (H_{u_R}^0 + v_u) + i H_{u_I}^0$$
(2.12)

$$H_d^0 = (H_{d_R}^0 + v_d) + i H_{d_I}^0$$
(2.13)

$$S = (S_R + v_s) + iS_I$$
 (2.14)

$$T^{0} = (T^{0}_{R} + v_{t}) + iT^{0}_{I}$$
(2.15)

The minimization conditions are derived from the fact that,

$$\frac{\partial V}{\partial v_u} = \frac{\partial V}{\partial v_d} = \frac{\partial V}{\partial v_s} = \frac{\partial V}{\partial v_t} = 0$$
(2.16)

We can determine the soft breaking mass parameters like $m_{H_u}^2$, $m_{H_d}^2$, m_T^2 and m_S^2 using the following minimization conditions,

$$m_{H_u}^2 = \cot \beta [A_{eff} - (\lambda^2 + \frac{\lambda_2^2}{2}) \frac{v^2}{2} \sin 2\beta + \lambda \lambda_4 v_t^2 - \sqrt{2} \lambda_2 \lambda_4 v_t v_s - \frac{\lambda_2}{\sqrt{2}} A_{\lambda_2} v_t] - \mu_{eff}^2 + \frac{1}{4} (g_1^2 + g_2^2) v^2 \cos 2\beta$$
(2.17)

$$m_{H_d}^2 = \tan \beta [A_{eff} - (\lambda^2 + \frac{\lambda_2}{2})\frac{v^2}{2}\sin 2\beta + \lambda\lambda_4 v_t^2 - \sqrt{2}\lambda_2\lambda_4 v_t v_s - \frac{\lambda_2}{\sqrt{2}}A_{\lambda_2} v_t] - \mu_{eff}^2 - \frac{1}{4}(g_1^2 + g_2^2)v^2\cos 2\beta$$
(2.18)

$$m_{S}^{2} = v^{2} \left[\frac{v_{t}}{\sqrt{2}v_{s}} \lambda \lambda_{2} + \lambda \lambda_{1} \sin 2\beta + \frac{1}{2v_{s}} \lambda A_{\lambda} \sin 2\beta - \lambda^{2} \right] - \left[2\lambda_{1}^{2}v_{s} + \lambda A_{\lambda_{1}} \right] v_{s}$$
$$-\lambda_{4} v_{t}^{2} \left[B_{\lambda}/v_{s} + 2\lambda_{1} + 4\lambda_{4} \right] - \sqrt{2}\lambda_{2}\lambda_{4} v_{u} v_{d} v_{t}/v_{s}$$
(2.19)

$$m_T^2 = \left[\frac{1}{\sqrt{2}}\lambda\lambda_2\frac{v_s}{v_t} - \frac{\lambda_2^2}{2} - \frac{\lambda_2}{2\sqrt{2}v_t}A_{\lambda_2}\sin 2\beta\right]v^2 - 2\lambda_4^2v_t^2 + 2\lambda\lambda_4v_uv_d$$
$$-\lambda_4v_s^2\left[2B_\lambda/v_s + 2\lambda_1 + 4\lambda_4\right] - \sqrt{2}\lambda_2\lambda_4v_uv_dv_s/v_t \qquad (2.20)$$

where,

$$A_{eff} = \lambda v_s [A_\lambda + \lambda_1 v_s] \tag{2.21}$$

and $v_u^2 + v_d^2 = v^2 = (174)^2 GeV^2$, $\tan \beta = v_u/v_d$.

Due to the addition of the triplets, the gauge bosons receive additional contribution in their masses like,

$$M_Z^2 = \frac{1}{2}(g_1^2 + g_2^2)v^2$$
(2.22)

$$M_W^2 = \frac{1}{2}g_2^2(v^2 + 4v_t^2)$$
(2.23)

The ρ -parameter at the tree-level is defined as,

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2 \theta_W} = 1 + 4 \frac{v_t^2}{v^2}$$
(2.24)

Clearly, the ρ -parameter deviates from unity by a factor of $4\frac{v_t^2}{v^2}$. Using the recent bound on ρ -parameter at 95% C.L. we can determine the bound on the triplet Higgs vev v_t . ρ can be confined in the range 0.9799-1.0066 [101] and hence $v_t \leq 9$ GeV at 95% C.L.

2.3 Higgs Sector

2.3.1 CP-even Higgs Mass Matrices

The CP-even sector of this model has been extended with respect to the MSSM by an additional neutral singlet and triplet scalar fields, which will contribute significantly to the mass(es) of the CP-even Higgs(es). The symmetric CP-even Higgs mass matrix written in the gauge basis of ($H_{u_R}^0$, $H_{d_R}^0$, T_R^0 , S_R) has 10 independent components. After Electroweak symmetry breaking (EWSB) the entries of the squared mass-matrix are,

$$\begin{split} M_{11}^{2} &= \frac{1}{2}(g_{1}^{2} + g_{2}^{2})v^{2}\sin^{2}\beta + C_{1}\cot\beta + C_{4}, \\ M_{22}^{2} &= \frac{1}{2}(g_{1}^{2} + g_{2}^{2})v^{2}\cos^{2}\beta + C_{1}\tan\beta + C_{4}, \\ M_{33}^{2} &= 4\lambda_{4}^{2}v_{t}^{2} + \lambda_{2}v^{2}[\lambda v_{s} - (A_{\lambda_{2}} + 2\lambda_{4}v_{s})\sin\beta\cos\beta]/\sqrt{2}v_{t}, \\ M_{44}^{2} &= \lambda_{1}v_{s}[A_{\lambda_{1}} + 4\lambda_{1}v_{s}] + [v_{t}(\lambda\lambda_{2}\frac{v^{2}}{\sqrt{2}} - \lambda_{4}B_{\lambda}v_{t}) \\ &+ (\lambda A_{\lambda} - \sqrt{2}\lambda_{2}\lambda_{4}v_{t})v^{2}\sin\beta\cos\beta]/v_{s}, \\ M_{12}^{2} &= -C_{1} + [2\lambda^{2} + \lambda_{2}^{2} - \frac{(g_{1}^{2} + g_{2}^{2})}{2}]v^{2}\sin\beta\cos\beta, \\ M_{13}^{2} &= v[C_{2}\cos\beta - \sqrt{2}\lambda_{2}\mu_{eff}\sin\beta], \\ M_{14}^{2} &= -v[C_{3}\cos\beta - 2\lambda\mu_{eff}\sin\beta], \\ M_{23}^{2} &= v[C_{2}\sin\beta - \sqrt{2}\lambda_{2}\mu_{eff}\cos\beta], \\ M_{24}^{2} &= -v[C_{3}\sin\beta - 2\lambda\mu_{eff}\cos\beta], \\ M_{34}^{2} &= 2\lambda_{4}v_{t}[B_{\lambda} + 2v_{s}(\lambda_{1} + 2\lambda_{4})] - \lambda_{2}v^{2}(\lambda - 2\lambda_{4}\sin\beta\cos\beta)/\sqrt{2} \ (2.25) \end{split}$$

where C_i 's are defined as,

$$C_{1} = A_{eff} + \lambda \lambda_{4} v_{t}^{2} - \lambda_{2} A_{\lambda_{2}} \frac{v_{t}}{\sqrt{2}} - \sqrt{2} \lambda_{2} \lambda_{4} v_{t} v_{s},$$

$$C_{2} = \frac{\lambda_{2} A_{\lambda_{2}}}{\sqrt{2}} - 2\lambda \lambda_{4} v_{t} + \sqrt{2} \lambda_{4} \lambda_{2} v_{s},$$

$$C_{3} = \lambda A_{\lambda} + 2\lambda \lambda_{1} v_{s} - \sqrt{2} \lambda_{2} \lambda_{4} v_{t},$$

$$C_{4} = \lambda_{2} v_{t} [\frac{\lambda_{2} v_{t}}{2} - \sqrt{2} \lambda v_{s}] \qquad (2.26)$$

and A_{eff} is defined in equation(2.21). Therefore, the CP-even higgs sector consists of four massive higgs as h, H_1 , H_2 and H_3 .

• Bound on the lightest Higgs mass :

The bound on the lightest Higgs mass is derived from the fact that, the smallest eigenvalue of a real, symmetric $n \times n$ matrix is smaller than the smallest eigenvalue of the upper left 2×2 sub-matrix [100]. Using this we obtain an upper bound on the lightest CP-even Higgs mass,

$$m_h^2 \leqslant M_Z^2 \left[\cos^2 2\beta + \frac{2\lambda^2}{g_1^2 + g_2^2} \sin^2 2\beta + \frac{\lambda_2^2}{g_1^2 + g_2^2} \sin^2 2\beta \right]$$
(2.27)

The bound on lightest Higgs mass has been considerably improved over the MSSM due to the additional contribution from the singlet and triplet gauge fields. Using equation(2.27) we can put constraints on the parameters like λ , λ_2 and tan β satisfying the recent bound on Higgs mass from ATLAS and CMS.

2.3.2 CP-odd Higgs Mass Matrices

The elements of the 4×4 CP-odd Higgs squared mass matrix, after EWSB, in the basis of ($H_{d_I}^0$, $H_{u_I}^0$, S_I , T_I^0) are,

$$\begin{split} M_{P_{11}}^2 &= C_1 \tan \beta + C_4, \\ M_{P_{22}}^2 &= C_1 \cot \beta + C_4, \\ M_{P_{33}}^2 &= -3\lambda_1 A_{\lambda_1} v_s - \lambda_4 [B_{\lambda} + 4\lambda_1 v_s] \frac{v_t^2}{v_s} + D_1 (\frac{v_t}{v_s}) + [\lambda A_{\lambda}/v_s + 4\lambda\lambda_1] v^2 \sin \beta \cos \beta, \\ M_{P_{44}}^2 &= -4\lambda_4 v_s [B_{\lambda} + \lambda_1 v_s] + D_1 (\frac{v_s}{v_t}) + [4\lambda\lambda_4 - \frac{1}{\sqrt{2}v_t}\lambda_2 A_{\lambda_2}] v^2 \sin \beta \cos \beta, \\ M_{P_{12}}^2 &= A_{eff} - \frac{v_t}{\sqrt{2}}\lambda_2 A_{\lambda_2} + \lambda_4 v_t [\lambda v_t - \sqrt{2}\lambda_2 v_s], \\ M_{P_{13}}^2 &= v \sin \beta [\lambda A_{\lambda} - 2\lambda\lambda_1 v_s + \sqrt{2}\lambda_2 \lambda_4 v_t], \\ M_{P_{14}}^2 &= -v \sin \beta [2\lambda\lambda_4 v_t + \frac{1}{\sqrt{2}}\lambda_2 (A_{\lambda_2} - 2\lambda_4 v_s)], \\ M_{P_{23}}^2 &= M_{P_{13}}^2 / \tan \beta, \\ M_{P_{24}}^2 &= M_{P_{14}}^2 / \tan \beta, \\ M_{P_{34}}^2 &= -2\lambda_4 v_t (B_{\lambda} - 2\lambda_1 v_s) - D_1 \end{split}$$

$$(2.28)$$

where,

$$D_1 = \frac{1}{\sqrt{2}}\lambda_2 v^2 (\lambda + 2\lambda_4 \sin\beta\cos\beta)$$
 (2.29)

Thus, the CP-odd higgs sector contains three pseudo-scalar Higgs A_1 , A_2 and A_3 . This matrix always contains a Goldstone mode G^0 (gives mass to Z-boson), which can be written as,

$$G^{0} = \cos\beta H^{0}_{d_{I}} - \sin\beta H^{0}_{u_{I}} \tag{2.30}$$

and we rotate the mass matrix in the basis (G^0, A_1, A_2, A_3) where,

$$\begin{pmatrix} A_1 \\ G^0 \\ A_2 \\ A_3 \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta & 0 & 0 \\ -\sin\beta & \cos\beta & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} H_{u_I}^0 \\ H_{d_I}^0 \\ S_I \\ T_I^0 \end{pmatrix}$$
(2.31)

After removing the Goldstone mode, we again rotate the remaining 3×3 mass matrix and finally obtain,

$$P_{1} = \cos \alpha \sin \beta H_{d_{I}} + \cos \alpha \cos \beta H_{u_{I}} + \sin \alpha S_{I},$$

$$P_{1} = -\sin \alpha \sin \beta H_{d_{I}} - \sin \alpha \cos \beta H_{u_{I}} + \cos \alpha S_{I},$$

$$P_{3} = T_{I} \qquad (2.32)$$

where P_1 , P_2 , P_3 are the massive modes.

2.3.3 Charged Higgs Mass Matrices

The charged Higgs sector comprises of a 4×4 symmetric matrix, written in the basis $(H_u^+, H_d^{-*}, T_0^+, T_0^{-*})$, which has 10 independent components, (after EWSB) given by

$$(M_{\pm}^{2})_{11} = E_{1}v_{d}^{2} + [\sqrt{2}\lambda\lambda_{2}v_{t}v_{s} + \frac{\lambda_{2}^{2}}{2}v_{t}^{2}] + C_{1}\cot\beta,
(M_{\pm}^{2})_{12} = A_{eff} + E_{1}v_{u}v_{d} + [\lambda_{2}A_{\lambda_{2}} + \sqrt{2}\lambda v_{t} + 2\lambda_{2}v_{s}]\frac{v_{t}}{\sqrt{2}},
(M_{\pm}^{2})_{13} = E_{2}v_{d} - 2\lambda_{2}v_{u}[\lambda v_{s} + \frac{\lambda_{2}v_{t}}{\sqrt{2}}],
(M_{\pm}^{2})_{14} = E_{3}v_{d} + v_{u}\lambda_{2}\mu_{eff},
(M_{\pm}^{2})_{22} = E_{1}v_{u}^{2} + [\sqrt{2}\lambda\lambda_{2}v_{t}v_{s} + \frac{\lambda_{2}^{2}}{2}v_{t}^{2}] + C_{1}\tan\beta,
(M_{\pm}^{2})_{23} = E_{3}v_{u} + v_{d}\lambda_{2}\mu_{eff},
(M_{\pm}^{2})_{24} = E_{2}v_{u} - 2\lambda_{2}v_{d}[\lambda v_{s} + \frac{\lambda_{2}v_{t}}{\sqrt{2}}],
(M_{\pm}^{2})_{33} = \frac{g_{2}^{2}}{2}[v_{u}^{2} - v_{d}^{2}] + \lambda_{2}^{2}v_{u}^{2} + E_{4},
(M_{\pm}^{2})_{34} = [g_{2}^{2} - 2\lambda_{4}^{2}]v_{t}^{2} - 2\lambda_{4}v_{s}[B_{\lambda} + \lambda_{1}v_{s}] + 2\lambda\lambda_{4}v_{u}v_{d},
(M_{\pm}^{2})_{44} = \frac{g_{2}^{2}}{2}[v_{d}^{2} - v_{u}^{2}] + \lambda_{2}^{2}v_{d}^{2} + E_{4}$$

$$(2.33)$$

where E_i 's are defined as,

$$E_{1} = \frac{g_{2}^{2}}{2} - \lambda^{2} + \frac{\lambda_{2}^{2}}{2},$$

$$E_{2} = \frac{g_{2}^{2}v_{t}}{\sqrt{2}} + 2\lambda_{2}\lambda_{4}v_{s},$$

$$E_{3} = \frac{g_{2}^{2}v_{t}}{\sqrt{2}} - \lambda_{2}A_{\lambda_{2}},$$

$$E_{4} = g_{2}^{2}v_{t}^{2} + 4\lambda_{4}^{2}v_{s}^{2}$$
(2.34)

After diagonalization, we obtain one massless Goldstone state G^+ (gives mass to W^{\pm} -boson, since $G^- \equiv G^{+*}$),

$$G^{+} = \sin\beta H_{u}^{+} - \cos\beta H_{d}^{-*} + \sqrt{2}\frac{v_{t}}{v}(T_{0}^{+} - T_{0}^{-*})$$
(2.35)

and three other massive modes like $H_1^\pm, H_2^\pm, H_3^\pm$

2.4 Neutralinos and Charginos

In the fermionic sector, the neutral component of the triplet and singlet i.e, \tilde{T}^0 and \tilde{S} mix with the higgsinos and the gauginos. Thus, the neutralino mass matrix extended by the singlet and triplet sector, in the basis $(\tilde{B}, \tilde{W}^0, \tilde{H}^0_d, \tilde{H}^0_u, \tilde{S}, \tilde{T}^0)$ is given by,

$$\mathcal{M}_{\bar{G}} = \begin{pmatrix} M_1 & 0 & -c_\beta s_w M_Z & s_\beta s_w M_Z & 0 & 0\\ 0 & M_2 & c_\beta c_w M_Z & -s_\beta c_w M_Z & 0 & 0\\ -c_\beta s_w M_Z & c_\beta c_w M_Z & 0 & -\mu_{eff} & -\lambda v_u & \frac{\lambda_2}{\sqrt{2}} v_u\\ s_\beta s_w M_Z & -s_\beta c_w M_Z & -\mu_{eff} & 0 & -\lambda v_d & \frac{\lambda_2}{\sqrt{2}} v_d\\ 0 & 0 & -\lambda v_u & -\lambda v_d & 2\lambda_1 v_s & 2\lambda_4 v_t\\ 0 & 0 & \frac{\lambda_2}{\sqrt{2}} v_u & \frac{\lambda_2}{\sqrt{2}} v_d & 2\lambda_4 v_t & 2\lambda_4 v_s \end{pmatrix}$$
(2.36)

where, M_1 , M_2 are the soft breaking mass parameters for Bino and Wino respectively and

$$c_{\beta} = \cos \beta, \, s_{\beta} = \sin \beta, \, c_w = \cos \theta_w \text{ and } s_w = \sin \theta_w$$

The left-most 4×4 entries are exactly identical with that in MSSM, except the μ_{eff} -term which is defined in equation(2.5). As the triplet and the singlet fermion

does not have any interaction with the neutral gauginos the right-most 2×2 entries are zero.

The chargino mass terms in the Lagrangian can be written as,

$$-\frac{1}{2}[\tilde{G}^{+T}M_c^T.\tilde{G}^- + \tilde{G}^{-T}M_c.\tilde{G}^+]$$
(2.37)

where, the basis \tilde{G}^+ and \tilde{G}^- are specified as,

$$\tilde{G}^{+} = \begin{pmatrix} \tilde{W}^{+} \\ \tilde{H}_{u}^{+} \\ \tilde{T}^{+} \end{pmatrix}, \quad \tilde{G}^{-} = \begin{pmatrix} \tilde{W}^{-} \\ \tilde{H}_{d}^{-} \\ \tilde{T}^{-} \end{pmatrix}$$

Similarly, the charged component of the triplet, \tilde{T}^+ and \tilde{T}^- contribute to the chargino mass matrix. The chargino matrix in the gauge basis \tilde{G}^+ and \tilde{G}^- is given by,

$$\mathcal{M}_{ch} = \begin{pmatrix} M_2 & \frac{1}{\sqrt{2}}g_2v_d & g_2v_t \\ \frac{1}{\sqrt{2}}g_2v_u & \lambda v_s + \frac{\lambda_2}{\sqrt{2}}v_t & \lambda_2v_d \\ -g_2v_t & \lambda_2v_u & 2\lambda_4v_s \end{pmatrix}$$
(2.38)

where,

$$\tilde{G}^{+} = \begin{pmatrix} \tilde{W}^{+} \\ \tilde{H}_{u}^{+} \\ \tilde{T}^{+} \end{pmatrix}, \quad \tilde{G}^{-} = \begin{pmatrix} \tilde{W}^{-} \\ \tilde{H}_{d}^{-} \\ \tilde{T}^{-} \end{pmatrix}$$

Since, $\mathcal{M}_{ch}^T \neq \mathcal{M}_{ch}$, this matrix is diagonalised via bi-unitary transformation, which requires two distinct unitary matrices U and V such that,

$$\tilde{\chi}^+ = V\tilde{G}^+,$$

 $\tilde{\chi}^- = U\tilde{G}^-$
(2.39)

The diagonal matrix reads,

$$U^* \mathcal{M}_{ch} V^{-1} = \begin{pmatrix} m_{\tilde{\chi}_1^{\pm}} & 0 & 0 \\ 0 & m_{\tilde{\chi}_2^{\pm}} & 0 \\ 0 & 0 & m_{\tilde{\chi}_3^{\pm}} \end{pmatrix}$$
(2.40)

and similarly the hermitian conjugate of eqn.2.40 also gives diagonal chargino mass matrix.

2.5 Results and Discussions

2.5.1 Choice of Benchmark Sets

The main results of this paper are shown in Table(2.1) and (2.2) [106]. We have specified the values of the parameters like couplings, soft-breaking parameters at the Electroweak (EW) scale. The choice of $\tan \beta$, λ and λ_2 are restricted from the bound on lightest Higgs mass(2.27). In FIG.2.1 we show relation between λ_2 and λ for different values of $\tan \beta$. As we increase $\tan \beta$, λ and λ_2 tend to shift towards the higher values.

Plot in the right-hand panel of FIG.2.1 shows the dependence of m_h on $\tan \beta$ for some particular choices of $\lambda = 0.6$, 0.64 and $\lambda_2 = 0.75$, 1.02, which are consistent with $m_h = 125$ GeV (shown in the dotted line). In order to satisfy the bound on Higgs mass, we can put constraint on $\tan \beta$ which is, $\tan \beta \leq 3.0$. The coupling λ_1 sets the mass for the singlino through the Yukawa term $2\lambda_1 S \chi_s \cdot \chi_s$. In order to have a light neutralino for satisfying the dark matter phenomenology we choose small values of $\lambda_1 = 0.2, 0.25$ as our benchmark values. The choice of λ_4 is determined from the bounds on chargino masses. The other soft breaking parameters $A_{\lambda}, A_{\lambda_1}, A_{\lambda_2}, B_{\lambda}$ are chosen to fit the CP-even scalar masses specially to make the lightest Higgs mass close to 125 GeV. Finally, we have chosen μ_{eff} to be $\mathcal{O}(200 \text{ GeV})$ and $v_t = 2 \text{ GeV}$, which determines the choice of v_s from equation.(2.5). The ratio of M_1 to M_2 at the electroweak scale is consistent with universal gaugino masses at GUT scale and gravity mediated SUSY breaking.

The mass spectrum shown in Table(2.2) indicates all masses at the tree-level. The Higgs spectrum consists of 4 CP-even Higgs (h, H_1 , H_2 , H_3), 3 pseudo-scalar Higgs (A_1 , A_2 , A_3) and 3 charged Higgs (H_1^{\pm} , H_2^{\pm} , H_3^{\pm}). We obtain significant contribution from the singlet and triplet sector at the tree-level which is highly appreciable, since this has raised the mass of the lightest CP-even Higgs boson to 125 GeV. Here we do not require a significant radiative contribution from the top-stop-sector [87].



Figure 2.1: Top : Plot of λ vs. λ_2 , for $\tan \beta = 2$ (dashed), 3 (thick) with $m_h = 125$ GeV. Bottom : Plot of m_h vs. $\tan \beta$ for $\lambda = 0.6, \lambda_2 = 0.75$ (dashed) $\lambda = 0.64, \lambda_2 = 1.02$ (thick) and the dotted line shows the recent bound i.e. $m_h = 125$ GeV

The components of the lightest physical Higgs for $\tan \beta = 2.0$ are given as,

$$h = 0.84205H_{u_R}^0 + 0.44422H_{d_R}^0 + 0.01977T_R^0 + 0.30533S_R$$
(2.41)

Parameters at EW scale	Point 1	Point 2
aneta	2.0	3.0
λ	0.60	0.64
λ_1	0.20	0.25
λ_2	0.75	1.02
λ_4	0.17	0.20
$\mu_{eff}[{ m GeV}]$	200	200
$A_{\lambda}[{ m GeV}]$	400	500
$A_{\lambda_1}[\text{GeV}]$	-10	-10
$A_{\lambda_2}[\text{GeV}]$	600	700
$B_{\lambda}[{ m GeV}]$	500	600
$v_t[\text{GeV}]$	2	2
$M_1[\text{GeV}]$	150	200
$M_2[\text{GeV}]$	300	400

Table 2.1: Value of the parameters specified at the Electroweak scale for two sets of Benchmark points.

The lightest Higgs mass eigenstate has significant contribution from the Singlet and some contribution from the Triplet sectors. We obtain the lightest scalar Higgs mass for the two sets of benchmark points as 120.6 GeV and 119.2 GeV respectively. This will change the $h \rightarrow \gamma \gamma$ branching compared to the standard model and precise determination of the Higgs decay branchings at LHC will be a good test of this model. In the pseudo-scalar Higgs sector, we obtain one Goldstone boson exactly identified as equation(2.30),i.e. $G^0 = 0.4472H_{d_I}^0 0.8942H_{u_I}^0$, for tan $\beta = 2.0$ and $G^0 = 0.3163H_{d_I}^0 - 0.9487H_{u_I}^0$, for tan $\beta = 3.0$. All other Higgs masses are listed in Higgs spectrum of Table(2.2).

The neutralino and the chargino sector consists of six and three mass-eigenstates respectively. The mass of the lightest neutralino being $\mathcal{O}(100 \text{ GeV})$, is the LSP of this model. The prospects of the LSP being a dark matter candidate is discussed in Section(2.5.2). Rest of the mass-spectrum are shown in Table(2.2)

Mass Spectrum	Point 1	Point 2		
Neutral Higgs Spectrum				
$m_h^{Tree}[\text{GeV}]$	120.6	119.2		
$m_{H_1}[\text{GeV}]$	145.5	156.8		
$m_{H_2}[\text{GeV}]$	482.4	630.7		
$m_{H_3}[{ m GeV}]$	825.2	707.9		
$m_{A_1}[\text{GeV}]$	114.3	116.9		
$m_{A_2}[\text{GeV}]$	487.8	629.9		
$m_{A_3}[\text{GeV}]$	897.3	816.0		
Charged Higgs Spectrum				
$m_{H_1}^{\pm}[\text{GeV}]$	208.4	239.9		
$m_{H_2}^{\pm}[\text{GeV}]$	280.5	320.6		
$m_{H_0}^{\pm}[\text{GeV}]$	496.3	647.1		
113 1				
Neutralino Spe	ectrum			
Neutralino Sp $m_{\tilde{\chi}_1^0}[\text{GeV}]$	e ctrum 100.4	102.9		
$m_{\tilde{\chi}_1^0}$ [GeV] $m_{\tilde{\chi}_2^0}$ [GeV]	ectrum 100.4 122.6	102.9 145.7		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7	102.9 145.7 205.9		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6	102.9 145.7 205.9 261.5		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$ $m_{\tilde{\chi}_5^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2	102.9 145.7 205.9 261.5 265.7		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$ $m_{\tilde{\chi}_5^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0	102.9 145.7 205.9 261.5 265.7 426.6		
$m_{\tilde{\chi}_{1}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{2}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{2}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{6}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{6}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{6}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{6}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{6}^{0}}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0 trum	102.9 145.7 205.9 261.5 265.7 426.6		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$ $m_{\tilde{\chi}_5^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0 trum 124.2	102.9 145.7 205.9 261.5 265.7 426.6 127.7		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$ $m_{\tilde{\chi}_5^0}[\text{GeV}]$ $m_{\tilde{\chi}_6^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_1^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0 trum 124.2 194.5	102.9 145.7 205.9 261.5 265.7 426.6 127.7 250.2		
$m_{\tilde{\chi}_1^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_2^0}[\text{GeV}]$ $m_{\tilde{\chi}_3^0}[\text{GeV}]$ $m_{\tilde{\chi}_4^0}[\text{GeV}]$ $m_{\tilde{\chi}_5^0}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0 trum 124.2 194.5 347.1	102.9 145.7 205.9 261.5 265.7 426.6 127.7 250.2 428.1		
$m_{\tilde{\chi}_{1}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{2}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{2}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{4}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{0}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{5}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{2}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{1}}[\text{GeV}]$ $m_{\tilde{\chi}_{3}^{1}}[\text{GeV}]$	ectrum 100.4 122.6 164.7 212.6 248.2 345.0 trum 124.2 194.5 347.1	102.9 145.7 205.9 261.5 265.7 426.6 127.7 250.2 428.1		

Table 2.2: Mass Spectrum and Relic Density for two sets of Benchmark points.

2.5.2 Dark Matter

We have analyzed the neutralino sector where the lightest neutralino (LSP), is a mixture of Higgsino-Triplino and turns out to be a viable Dark Matter candidate. The components of $\tilde{\chi_0}$ (for tan $\beta = 2.0$), i.e. the LSP are,

$$\tilde{\chi_0} = -0.321\tilde{B} + 0.192\tilde{W_3^0} - 0.323\tilde{H_d^0} + 0.644\tilde{H_u^0} - 0.213\tilde{S} + 0.544\tilde{T^0}$$
(2.42)

Since the LSP has mass $\mathcal{O}(100 \text{ GeV})$, there are two possibilities of final states into which it can annihilate, i.e. (i)Fermion final states and (ii)Gauge Boson final states. For annihilation into fermions, except $t\bar{t}$ it can go to any other $f\bar{f}$ pairs via pseudo-scalar Higgs, Z-boson exchange and sfermion exchange. But, if we consider the neutralino to be more like triplino, then its coupling with Z-boson is forbidden. Generally, it can annihilate into gauge boson pairs via several processes like chargino exchange, scalar Higgs exchange and Z-boson exchange. But the dominant contribution comes from annihilation into W^{\pm} via chargino exchange, which finally leads to the Relic Density of 0.117, consistent with WMAP [104].

The scalar interaction between the dark matter (i.e Neutralino LSP) and the quark is given by,

$$\mathcal{L}_{scalar} = a_q \bar{\chi} \chi \bar{q} q \tag{2.43}$$

where a_q is the coupling between the quark and the Neutralino. The scalar cross section for the Neutralino scattering off a target nucleus (one has to sum over the proton and neutrons in the target) is given by,

$$\sigma_{scalar} = \frac{4m_r^2}{\pi} (Zf_p + (A - Z)f_n)^2$$
(2.44)

where, m_r is the reduced mass of the nucleon and $f_{p,n}$ is the Neutralino coupling to proton or neutron[22, 29], given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q}, \qquad (2.45)$$

where $f_{Tu}^{(p)} = 0.020 \pm 0.004, f_{Td}^{(p)} = 0.026 \pm 0.005, f_{Ts}^{(p)} = 0.118 \pm 0.062, f_{Tu}^{(n)} = 0.014 \pm 0.003, f_{Td}^{(n)} = 0.036 \pm 0.008$ and $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [107]. $f_{TG}^{(p,n)}$ is related

to these values by

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}.$$
(2.46)

The term in Eq. 2.45 which includes $f_{TG}^{(p,n)}$ results from the coupling of the WIMP to gluons in the target nuclei through a heavy quark loop.

We can approximate $\frac{a_q}{m_q} \simeq \alpha/(s - m_h^2)$ where, α is the product of different coupling and mixings, m_q is the mass of the quark and $s = 4m_{\chi}^2$ (m_{χ} being the Dark matter mass). The parameter α plays a crucial role in determining the spin-independent cross-section and is highly model dependent. Using this we estimate $\alpha \simeq 2 \times 10^{-4} GeV^{-1}$ and the value of the spin-independent cross-section is $10^{-43} cm^2$, which is below the exclusion limits of XENON100 [105] and other direct detection experiments.

2.6 One-loop Correction to the Lightest Physical Higgs Mass

The one-loop correction to m_h^2 is calculated by constructing the Coleman-Weinberg potential [108],

$$V_{CW} = \frac{1}{64\pi^2} STr[M^4 (ln\frac{M^2}{Q_r^2} - \frac{3}{2})]$$
(2.47)

where M^2 are the field dependent tree-level mass matrices and Q_r is the renormalization scale. STr is the supertrace which includes a factor of $(-1)^{2J}(2J+1)$ and summed over the spin degrees of freedom. The one loop mass matrix can be derived from the above potential as follows,

$$(\Delta M_f^2)_{ij} = \frac{\partial^2 V_{CW}(f)}{\partial f_i \partial f_j}|_{vev} - \frac{\delta_{ij}}{\langle f_i \rangle} \frac{\partial V_{CW}(f)}{\partial f_i}|_{vev}$$
(2.48)

where, $f_{i,j}$ stands for all the real components of H_u^0 , H_d^0 , S and T^0 . Finally, the set of mass eigenvalues of the CP-even, CP-odd, Charged Higgs and Neutralino-Chargino mass matrices (all field-dependent) enters the calculation. The dominant contribution in the one-loop correction comes from the top-stop sector and the triplet sector. We compute the corrections only numerically using the bench-

Benchmark Point	m_h^{Tree} [GeV]	$m_h^{Tree+Loop}$ [GeV]
Point 1	120.6	124.9
Point 2	119.2	125.5

mark values assigned for the sets of parameters. The results we obtain [106] are given in Table(2.3)

Table 2.3: Value of the lightest physical Higgs mass after 1-loop correction for two sets of Benchmark points.

In both the cases we do not require large contribution from the radiative corrections to raise the lightest physical Higgs mass so as to satisfy the value of 125 GeV. This in turn implies that the contribution from the stop-top sector is not significant as in the case of MSSM. In fact in absence of fine tuning the correction to lightest physical Higgs mass from the stop-top sector is given by,

$$\delta m_{H_u}^2(Q) \simeq \frac{3m_t^2}{(4\pi)^2 v^2} ln \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2}$$
(2.49)

For, $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$ being \mathcal{O} (200 GeV), this amounts to a correction of only a few GeV.

2.7 Fine Tuning in the Electroweak Sector

In this model, the lightest physical Higgs mass at the tree-level is boosted compared to NMSSM, and other triplet extended model [102] as it gets contribution from both singlet and triplet sector(2.27). Therefore, we can obtain a Higgs boson close to 125 GeV even at the tree level. After including the leading order radiative corrections from the stop-top and triplet sector, we get

$$\delta m_{H_u}^2(Q) \simeq \frac{3y_t^2}{8\pi_2} (m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2 + A_t^2) ln(\frac{Q}{M_z}) + \frac{3\lambda_2^2}{8\pi_2} (m_T^2 + A_{\lambda_2}^2) ln(\frac{Q}{M_z})$$
(2.50)

where, $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$ are the soft masses of the stops, A_t is the soft trilinear coupling, y_t is the Yukawa coupling and Q is the fundamental scale of SUSY-breaking.

The fine-tuning parameter can be quantified [109, 110]as,

$$\Delta_{FT} \equiv \frac{m_{H_u}^2}{M_z^2} \frac{\partial M_z^2}{\partial m_{H_u}^2} \tag{2.51}$$

In case of MSSM (only first term in eqn.2.50 is present), we have

$$\Delta_{FT}^{Stop} \simeq \frac{3y_t^2}{8\pi_2} (m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2 + A_t^2) ln(\frac{Q}{M_z})$$
(2.52)

But the tree-level bound on Higgs mass is $m_h \leq M_z \cos 2\beta$. Therefore, one is forced to consider large values for $m_{\tilde{t}_1}$, $m_{\tilde{t}_2}$ and A_t , say 1 TeV in order to raise the lightest physical Higgs Boson mass upto 125 GeV. In this case, $\Delta_{FT}^{Stop} \simeq 80$ and thus it leads to maximal stop mixing.

In NMSSM, the radiative correction needed to get a 125 GeV Higgs mass is $\delta m_h = 55$ GeV. There is no doubt an improvement over MSSM, but still fine tuning is required in the stop-top sector[87]. In the model with one triplet [102], the lightest physical Higgs mass can be raised to 113 GeV. Here, the required value of radiative correction is $\delta m_h = 53$ GeV. Now, the fine tuning due to the triplet sector is,

$$\Delta_{FT}^{Trip} \simeq \frac{3\lambda_2^2}{8\pi_2} (m_T^2 + A_{\lambda_2}^2) ln(\frac{Q}{M_z})$$
(2.53)

where $\lambda_2 = 0.8, 0.9$. The value of Δ_{FT}^{Trip} can be as large as 40. Therefore, this model can no longer be considered as a zero fine-tuning model.

Now coming to our model, we require $\delta m_h \simeq 35$ GeV only- here we see a distinct improvement of 20-50 GeV compared to other models discussed so far. Also, $\lambda_2 = 0.75, 1.02$ being comparable to y_t , we do not need heavy stops or large stop-top mixing to get the required Higgs mass. For example, $m_T = 200$ GeV, $A_{\lambda_2} = 700$ GeV and Q = 1 TeV, we obtain $\Delta_{FT}^{Trip} \simeq 10$ [106]. Thus, we can achieve little fine-tuning compared to other models, since the lightest physical Higgs mass can be large at tree level and does not require large contribution from the radiative corrections. Here we note that, the Higgs-Triplet-Higgs coupling λ_2 (0.75 and 1.02) becomes non-perturbative at GUT scale. But, these choices of λ_2 actually helps to raise the Higgs mass close to 125 GeV at the tree-level. Other alternative corrections. Therefore, we improve the level of fine-tuning at the cost of giving up perturbativity of λ_2 at GUT scale.

2.8 Conclusion and Outlook

In this chapter, we have explored an extension of MSSM where the Higgs sector is extended by a singlet and a Y = 0 triplet superfield. This is the minimal model which gives a tree level Higgs mass of $\mathcal{O}(119\text{-}120 \text{ GeV})$ and the one-loop correction can easily raise it to 125 GeV without significant contribution from the stop-top sector. However, $\lambda_2 = 0.75, 1.02$ (at Electroweak scale) becomes nonperturbative at the GUT scale, while all other couplings remain perturbative upto GUT scale - on the other hand this is the price we pay to retain small fine-tuning.

In addition, we see that the triplino and singlino contributions to the neutralino mass matrix gives a viable dark matter candidate with mass around 100 GeV which may be seen at the LHC from the missing transverse energy signals [41, 111]. In MSSM and NMSSM the problem for getting the correct relic density of dark matter is related to the necessity of choosing chargino and scalar masses to be in the multi TeV scale to fit the Higgs mass from radiative corrections. The DM mass in MSSM is around 700 GeV while in NMSSM it is possible to obtain viable DM in the 100 GeV range. The main advantage of our model for the dark matter is that since the sparticle masses need not be very large compared to the electroweak scale the 'WIMP miracle' is restored and we are able to get DM mass in the 100 GeV range over a large parameter space of our model.

The data from LHC with integrated luminosity of $5fb^{-1}$ has not only given an indication the Higgs mass but there is also a measurement of the Higgs decay branchings into different channels. Detailed analysis [112] of the 125 GeV Higgs branching fractions seen at the LHC indicates that the signal ratio for Higgs decay into two photons is larger than SM prediction by a factor of 2.0 ± 0.5 , decay into WW^* and ZZ^* channels is smaller than SM by a factor of 0.5 ± 0.3 and into bb and $\tau\tau$ channels it is factor 1.3 ± 0.5 consistent with SM. The lightest CP-even Higgs (2.41) has a sizable fraction of the singlet and the Higgs decay phenomenology will be distinguishable from the MSSM [113] and likely to be similar to the NMSSM scenario [114–116]. But, there will be some contribution from the triplet sector too. The phenomenological aspects of the real-triplet extended SM has been studied in [117]. More data from LHC will pinpoint or rule out the extended Higgs sector models and it would be useful to study the singlet-triplet extended MSSM model in greater detail with emphasis on the LHC signal in the future.

Chapter 3

Explaining 130 GeV monochromatic γ -ray line in TSMSSM

3.1 Overview

Recently it has been pointed out [118–121] that the analysis of the Fermi-LAT gamma-ray data [68] reveals the existence of a peak at around 130 GeV coming from the vicinity of the galactic center. Further, it shows that the interpretation of the gamma ray peak as due to DM annihilation with mass $129.8 \pm 2.4^{+7}_{-13}$ GeV and annihilation cross-section $\langle \sigma v \rangle_{\gamma\gamma} = (1.27 \pm 0.32^{+0.18}_{-0.28}) \times 10^{-27} \text{cm}^3 \text{sec}^{-1}$ fits the signal well. Numerous studies have been made to accommodate this feature in terms of DM annhibition in both model-independent way [122] and in specificifically Standard Model (SM) extended by singlets and triplet [123–125]. After the discovery of the Higgs-like boson around mass window 125-126 GeV, there is another intriguing possibility of a signal beyond SM in the $h\,\rightarrow\,\gamma\gamma$ channel. The ratio between the Higgs di-photon decay rate observed at LHC and the one expected in the SM is $R_{\gamma\gamma} = 1.65^{+0.34}_{-0.30}$ for ATLAS ($m_h = 126 \text{ GeV}$) whereas CMS have now fallen down to $R_{\gamma\gamma} = 0.78^{+0.28}_{-0.26}$ for $m_h = 125 \text{ GeV}$ [126]. This channel will be an important discriminator of models as future LHC data pinpoints this number more precisely. The implications of the modified diphoton decay width in a generic model independent approach have been discussed in ref. [113]. Very recently, a vector Higgs-portal dark matter model (SM extended by $U(1)_x$ gauge symmetry) [127] has addressed both Fermi-gamma ray line and diphoton excess simultaneously.

It is well-known that SUSY is the simplest model from protecting the Higgs mass from large radiative corrections without fine tuning. In the minimal supersymmetric standard model (MSSM) [35–40], the Higgs mass is close to the Z-boson mass at the tree level, which demands a large radiative correction to raise the Higgs mass to 125-126 GeV seen at the LHC [1, 2]. This in turn pushes the squark masses in the TeV range and hence the mixing in the top-stop sector becomes significant. This raises issues about fine tuning - which is somewhat solved by the so-called Next-to-minimal supersymmetric standard model (NMSSM) by adding a singlet chiral superfield to MSSM [83–86]. But to achieve a tree level Higgs mass close to 125 GeV, we need a large $\lambda SH_u H_d$ coupling which borders in the nonperturbative regime of λ [87]. Another popular extension is the triplet-extended MSSM models with a Y = 0, SU(2) triplet superfield [100-102], where the tree level contribution to the Higgs mass comes from the $\lambda_2 H_d T_0 H_d$ term. But, [102] shows that the tree-level Higgs mass can be raised atmost to 113 GeV, which would still require substantial loop corrections from stops. Other possibilities include models with two real triplets $(Y = \pm 1)$ and one singlet [103] - studied with a motivation to solve the μ -problem as well as to obtain a large correction to the lightest Higgs mass. But, the analysis of the fermionic sector as well as the dark matter of this model is cumbersome.

In MSSM, the neutralino LSP, being the favourite candidate for DM, annihilates into two photons via loop-suppressed processes [128, 129] - the crosssection for which is usually too small to explain the signal. But, with a bino-like LSP [130] and through the exchange of light slepton and sneutrino the observed $\sigma v_{\gamma\gamma}$ is achieved in MSSM. An alternate possibility is to incorporate the internal Bremsstrahlung (IB), which can also give sharp spectral features in the γ -ray spectrum [131–133]. In bino DM annihilation to final state fermions, the fermion mass suppression in the cross section is avoided if there is a final state photon with the fermion pair [128, 129]. In ref.[134] it was pointed out that a significant higgsino component in the DM would lead to a continuum gamma ray spectrum from W^{\pm} final states and would not be able to explain the gamma ray peak. To avoid this, IB from bino dominated LSP's is more promising but there is a problem in getting a natural SUSY model with 130 GeV bino DM which gives the correct the relic abundance. MSSM could accommodate the enhancement in the di-photon decay rate with highly mixed light staus and large tan β [135].

In addressing the problem of explaining the 130 GeV gamma ray features, NMSSM models are most widely studied [136–138]. In NMSSM, the neutralino DM(~ 130 GeV) annihilates into two photon via resonant channel through psedoscalar singlet Higgs ($m_{A_s} \sim 260$ GeV) and light charged particle loops. NMSSM can also successfully account for the excess seen in the $h \rightarrow \gamma \gamma$ channels [116, 139, 140], in the case of strong singlet-doublet mixing, although the partial width of $h \rightarrow b\bar{b}$ is highly reduced in these models. In a generalised version of NMSSM model(GNMSSM) [141] simultaneously both the signals from Fermi and LHC has been explained in the same benchmark scenario.

In Chap.2, it was shown by adding a hypercharge Y = 0, SU(2)-triplet and a singlet chiral superfield there is an extra tree-level contribution to the Higgs mass and it can be raised close to 125 GeV at the tree level. Hence, no large contributions from stop loops is needed to get the required Higgs mass which alleviates the fine tuning problem of fixing the stop mass to a high precision at the GUT scale. Therefore a significant improvement of the fine tuning is achieved with respect to MSSM, NMSSM and other triplet-SUSY models. In addition, the model contains a dark matter(DM) candidate of mass $\mathcal{O}(100)$ GeV, with a correct relic abundance. Enhancement of diphoton decay width has been studied well in the triplet extended SUSY models [102, 142–144], where the contributions from to the charginos and charged Higgs(triplet like, with large triplet coupling) are taken into account. But, so far no benchmark points have been found which at the same time provide a viable DM in triplet extended SUSY models.

In this chapter, we attempt to explain the 130 GeV monochromatic gamma ray spectral feature in the triplet-singlet extended MSSM, i.e TSMSSM [106] (which was introduced in Chap.2) through the resonant annihilation of neutralino LSP into photons via pseudoscalar triplet Higgs of mass ~ $2m_{DM}$, which couples to the DM via the Yukawa term, $\lambda_2 T_0 \tilde{H}_u^0 \tilde{H}_d^0$. In addition, our model predicts a second photon peak at around 114 GeV with the cross-section being 0.75 times $\langle \sigma v \rangle_{\gamma\gamma}$. This DM has a correct relic abundance of 0.109 where dominant contribution comes from $\langle \sigma v \rangle_{W^+W^-}$. The spin-independent direct detection cross-section is well-below the latest XENON100 [55, 56] exclusion limits. Another motivation of this work is to provide an enhanced diphoton decay rate compared to SM through the additional contribution from the light chargino loops. This would be a specific prediction of our model and can be tested in the future collider search.

This chapter is organised as follows: In section 3.2, we attempt to provide an explanation for the Fermi-LAT monochromatic gamma ray line features with a neutralino LSP pair annihilation into two photon via pseudoscalar Higgs triplet near resonance. We substantiate our claim with a specific benchmark scenario which satisfy all desired phenomenological requirements. We also discuss about the relic abundance and spin-independent scattering cross-section in section 3.2.1 and section 3.2.2 respectively. In section 3.3, we show a detail formulation of the diphoton Higgs decay width. We present a short summary and conclusions in the last section.

3.2 130 GeV Fermi gamma ray line in TSMSSM

In this TSMSSM model [106], the dark matter is the LSP $\tilde{\chi}_1^0$ which can be expressed in the gauge basis as,

$$\tilde{\chi}_1^0 = N_{11}\tilde{B} + N_{12}\tilde{W}_3^0 + N_{13}\tilde{H}_d^0 + N_{14}\tilde{H}_u^0 + N_{15}\tilde{S} + N_{16}\tilde{T}^0$$
(3.1)

where, N_{11}^2 is the bino-fraction, N_{12}^2 is the wino-fraction, $N_{13}^2 + N_{14}^2$ is the higgsino-fraction, N_{15}^2 and N_{16}^2 are the singlino and triplino-fraction respectively.

We scan the corresponding regions of the parameter space of the tripletsinglet model [106, 145] and tune the couplings and masses, such that they satisfy all desired phenomenological properties. The main results are shown in

Parameters at EW scale		
aneta	1.8	
λ	0.55	
λ_1	0.20	
λ_2	0.80	
λ_4	0.25	
$\mu_{eff}[\text{GeV}]$	246	
$A_{\lambda}[{ m GeV}]$	400	
$A_{\lambda_1}[\text{GeV}]$	-50	
$A_{\lambda_2}[\text{GeV}]$	297.6	
$B_{\lambda}[{ m GeV}]$	270	
$v_t [{\rm GeV}]$	2	
$M_1[\text{GeV}]$	154.5	
$M_2[\text{GeV}]$	375	

Table 3.1: A sample set of benchmark points for $\tan \beta = 1.8$ and $M_1 = 154.5$ GeV.

Table. (3.1,3.2,3.3) [145]. In Table.3.1, we show a sample set of benchmark points for a particular choice of tan $\beta = 1.8$ specifying all the parameters, couplings and soft masses at the EW scale.

- As shown in [106], the CP-even physical Higgs boson receives significant contribution from the singlet and triplet through the terms $\lambda \hat{S} \hat{H}_d \cdot \hat{H}_u$ and $\lambda_2 \hat{H}_d \cdot \hat{T}_0 \hat{H}_u$ and thus its mass is raised to 122.9 GeV at tree level (shown in Table.3.2). It requires a little contribution from the radiative corrections raise it to 126 GeV. This lightest CP-even Higgs is SM-like with large H_u^0 and H_d^0 component.
- A dominantly triplet-like pseudoscalar Higgs A_T with mass ~ 260.54 GeV can be obtained by adjusting the soft-trilinear couplings. The psedoscalar triplet A_T has no tree-level coupling with the SM fermions or Z-boson. It can interact with the neutralinos and charginos via the Yukawa term in

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Higgs Spectrum [GeV]		
m_h^{Tree}	122.93	
m_{H_1}	175.29	
m_{H_2}	457.27	
m_{H_3}	538.86	
m_{A_1}	142.12	
m_{A_2}	260.54	
m_{A_3}	534.56	
$m_{H_1}^{\pm}$	133.13	
$m_{H_2}^{\pm}$	365.61	
$m_{H_3}^{\pm}$	545.59	
Neutralino Masses [GeV]		
$m_{ ilde{\chi}_1^0}$	130.02	
$m_{ ilde{\chi}_2^0}$	189.0	
$m_{ ilde{\chi}^0_3}$	215.47	
$m_{ ilde{\chi}_4^0}$	269.30	
$m_{ ilde{\chi}_5^0}$	283.49	
$m_{ ilde{\chi}_6^0}$	414.20	
Chargino Masses [GeV]		
$m_{\tilde{\chi}_1^{\pm}}$	131.92	
$m_{\tilde{\chi}_2^{\pm}}$	299.38	
	422.24	

Table 3.2: Mass spectrum for the specific choice of benchmark points.

the lagrangian like $\lambda_2 A_T \tilde{H}_u^0 \tilde{H}_d^0$. Although the doublet-triplet mixing terms like $\frac{\lambda_2^2}{2} \left[|H_u^0|^2 + |H_d^0|^2 \right] |T^0|^2$ is present in the scalar potential, but A_T cannot decay into two CP-even Higgs boson, m_h . Therefore the width of A_T is small, i.e, $\Gamma_T \simeq 6.84$ MeV- which boosts the Breit-Weigner propagator and cross-section $\langle \sigma v \rangle_{\gamma\gamma}$.



Figure 3.1: The dominant diagram for the resonant pair annihilation of neutralino into two photons via psedoscalar triplet Higgs A_T

• The LSP $\tilde{\chi}_1^0$ is dominantly bino-like $(N_{11} \sim 0.84)$ but contains substantial higgsino-fraction $(N_{13} \sim -0.31 \text{ and } N_{14} \sim 0.36)$. By suitably tuning the soft masses M_1 and M_2 , the desired mass of 130 GeV is obtained. Varying M_1 between 150-160 GeV, we obtain $127 \leq M_{\tilde{\chi}^0}[\text{GeV}] \leq 133$. Here, μ eff~ 246 GeV being less than $v_s \sim 450$ GeV makes the singlino $(N_{15} \sim$ -0.19) and triplino-fraction $(N_{16} \sim 0.10)$ less in $\tilde{\chi}_1^0$. Again, since M_1 is lighter than μ -eff, we get a enhancement in the bino fraction compared to higgsino. But, the significant higgsino fraction is required to get large value of $\langle \sigma v \rangle_{\gamma\gamma}$ through the resonant annihilation via psedoscalar Higgs A_T and the light chargino loops. In FIG.3.1, the dominant diagram for the resonant pair annihilation of neutralino into two photons via psedoscalar triplet Higgs A_T is shown. The lightest chargino $\tilde{\chi}_1^+$ and the DM are almost degenarate and is also dominantly higgsino-like.

The pair annihilation of $\tilde{\chi}^0$, with mass $129.8 \pm 2.4^{+7}_{-13}$ GeV into two photon demands a cross-section of $\langle \sigma v \rangle_{\gamma\gamma} = (1.27 \pm 0.32^{+0.18}_{-0.28}) \times 10^{-27} cm^3 sec^{-1}$ in order to fit the Fermi-LAT signal [68].

A simplified form of the analytical expression of $\langle \sigma v \rangle_{\gamma\gamma}$ following [122],

$$\langle \sigma v \rangle_{\gamma\gamma} = \frac{\alpha^2 g_f^2 g_{\chi}^2}{256\pi^3} \frac{m_{\chi_1^+}^2}{[(4m_{DM}^2 - m_{A_T}^2)^2 + \Gamma_T^2 m_{A_T}^2]} \times [arctan[(m_{\chi_1^+}^2 - m_{DM}^2)/m_{DM}^2]^{-1/2}]^2$$
(3.2)



Figure 3.2: Plot of $\sigma v_{\gamma\gamma}$ as a function of psedoscalar mass M_{A_T} . The dashed line shows the maximum value of $\langle \sigma v \rangle_{\gamma\gamma} \simeq 1.249 \times 10^{-27} \text{cm}^3 \text{sec}^{-1}$.

where, g_{χ} and g_f are the couplings of psedoscalar Higgs A_T with DM and the charged fermion in the loop respectively. Here, we take the assumption that only the lightest chargino, with mass 131.9 GeV contributes significantly. Upto a crude approximation, $g_{\chi} \sim \lambda_2 N_{13} N_{14}$ and $g_f \sim \lambda_2 U_{12} V_{12}$, where U and V are diagonalising matrix for the charginos. Finally, in the resonance limit of $m_{A_T} \sim 2m_{DM}$ and $m_{\chi^+} \rightarrow m_{DM}$, the pair annihilation cross-section becomes $\sim 1.249 \times 10^{-27} cm^3 s^{-1}$. However, the mass of the triplet-like CP-odd scalar Higgs has to lie accidentally close to 260 GeV to a precision ≤ 1.5 GeV. FIG.3.2 shows the behaviour of $\sigma v_{\gamma\gamma}$ with the psedoscalar triplet mass near resonance, this clarifies the need of tuning of both $M_{\tilde{\chi}_1^0}$ and m_{A_T} [145].

• A second γ -ray line at 114 GeV : Apart from the monochromatic γ -ray line at 130 GeV, there is another intriguing hint for a second line at ~111 GeV [146, 147], where the best fit to the relative cross-section is $\langle \sigma v \rangle_{\gamma Z} / \langle \sigma v \rangle_{\gamma \gamma} = 0.66^{+0.71}_{-0.48}$ [121]. A second photon line at 114 GeV is expected from kinematics if there is a $Z\gamma$ final state in the annhibition of

Observables	
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to \gamma \gamma) \ [10^{-27} \text{cm}^3 \text{ s}^{-1}]$	1.249
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to Z \gamma) \ [10^{-27} \mathrm{cm}^3 \mathrm{~s}^{-1}]$	0.94
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to WW) \ [10^{-27} \mathrm{cm}^3 \mathrm{\ s}^{-1}]$	3.57
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to ZZ) \ [10^{-27} \text{cm}^3 \text{ s}^{-1}]$	0.62
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to b \bar{b}) \ [10^{-27} \text{cm}^3 \text{ s}^{-1}]$	0.045
$\langle \sigma v \rangle (\chi_1^0 \chi_1^0 \to \tau \bar{\tau}) \ [10^{-27} \text{cm}^3 \text{ s}^{-1}]$	0.082
Ωh^2	0.109
$\sigma(p)_{SI}[10^{-9} \mathrm{pb}]$	0.681
$R_{\gamma\gamma}$	1.24

Table 3.3: Value of different annihilation cross-sections, relic abundance, spinindependent scattering cross-section and di-photon decay rate

the pair of $\tilde{\chi}_1^0$,

$$E_{\gamma} = m_{\tilde{\chi}_{1}^{0}} \left(1 - \frac{m_{Z}^{2}}{4m_{\tilde{\chi}_{1}^{0}}^{2}}\right) \tag{3.3}$$

where, $E_{\gamma} = 114 \text{ GeV}$ for $m_{\tilde{\chi}_1^0} = 130 \text{ GeV}$. The cross-section for $\langle \sigma v \rangle_{\gamma Z}$ is calculated using an approximation of the formulae given in [148]. Here, we find that for the set of benchmark points presented in Table.3.3, $\langle \sigma v \rangle_{\gamma Z} \simeq$ $0.943 \times 10^{-27} \text{cm}^3 \text{s}^{-1}$.

3.2.1 Relic Abundance

Another issue with dark matter is to satisfy the correct relic abundance, which is difficult in case when it is dominantly higgsino-like since it couples to gauge boson very efficiently and thus leads to large pair annihilation cross-section. This kind of interaction can be reduced by an enhanced bino component. We find a neutralino DM with $N_{11} \sim 0.84, N_{13} \sim -0.31$ and $N_{14} \sim 0.36$, which makes the relic density 0.109. The pair annihilations into final state W^+W^- , ZZ, $b\bar{b}$, $\tau^+\tau^-$ are shown in Table.3.3, calculated using micrOMEGAs2.4 [149– 151]. Thus, a bino dominated but with a substantial higgsino component dark matter is preferable in order to satisfy the latest PLANCK result, i.e, $\Omega_{\chi}h^2 = 0.1199 \pm 0.0027$ at 68% CL [5] whereas the corresponding value from the 9-year WMAP data is $\Omega_{\chi}h^2 = 0.1148 \pm 0.0019$ [4].

3.2.2 Calculation of spin-independent cross-section

Starting from a low-energy neutralino-quark effective lagrangian for spin-independent interaction,

$$L_{eff} = a_q \bar{\tilde{\chi}}_1^0 \tilde{\chi}_1^0 \bar{q} q \tag{3.4}$$

where, a_q is the neutralino-quark coupling, we obtain the scattering cross section (spin-independent) for the dark matter off of a proton or neutron as,

$$\sigma_{scalar} = \frac{4m_r^2}{\pi} f_{p,n}^2 \tag{3.5}$$

where, m_r is the reduced mass of the nucleon and $f_{p,n}$ is the neutralino coupling to proton or neutron[22, 29], given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q},$$
(3.6)

where $f_{Tu}^{(p)} = 0.020 \pm 0.004, f_{Td}^{(p)} = 0.026 \pm 0.005, f_{Ts}^{(p)} = 0.118 \pm 0.062, f_{Tu}^{(n)} = 0.014 \pm 0.003, f_{Td}^{(n)} = 0.036 \pm 0.008$ and $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [107]. $f_{TG}^{(p,n)}$ is related to these values by

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}.$$
(3.7)

The term in Eq. 4.20 which includes $f_{TG}^{(p,n)}$ results from the coupling of the WIMP to gluons in the target nuclei through a heavy quark loop.

In deriving an approximate form of a_q/m_q we ignore contributions from the squark exchange diagrams because of the latest LHC bounds on squark masses [152, 153]. Thus, a_q receives significant contribution from the t-channel exchange of CP-even Higgs bosons. The analytical form of a_q goes roughly as,

$$\frac{a_q}{m_q} \simeq \frac{S_{\chi\chi h_i}}{m_{h_i}^2} S_{h_i q q} \tag{3.8}$$

where, $S_{\chi\chi h_i}$ is the coupling between the neutralino and the CP-even Higgs bosons. For, up-type quarks, $S_{h_i u u} = \frac{g_2}{2M_w \sin\beta} S_{i1}$ and down-type, $S_{h_i d d} = \frac{g_2}{2M_w \cos\beta} S_{i2}$. Now, the coupling $S_{\chi\chi h_i}$ is a product of different combinations of λ 's, N_{1k} and $S_{i1,2}$. S_{ij} is the matrix which diagonalises the CP-even Higgs matrix, and the weak eigenstate basis is $(H_{u_R}^0, H_{d_R}^0, T_R^0, S_R)$. N_{1k} 's are the different components of the lightest neutralino dark matter. Under the assumption that only the lightest physical Higgs boson, i.e, h_1 ($m_{h_1} \simeq 125.8$ GeV) contributes dominantly, $S_{\chi\chi h_1}$ takes the form [145],

$$S_{\chi\chi h_{1}} \simeq g_{2}(N_{12} - \tan \theta_{W} N_{11})(S_{11}N_{13} - S_{12}N_{14}) -\sqrt{2}\lambda(S_{11}N_{14}N_{15} + S_{12}N_{13}N_{15} + S_{14}N_{14}N_{13}) + \sqrt{2}\lambda_{1}S_{14}N_{15}^{2} +\lambda_{2}(S_{11}N_{16}N_{13} + S_{12}N_{16}N_{14} + S_{13}N_{13}N_{14}) +\sqrt{2}\lambda_{4}(S_{14}N_{16}^{2} + 2S_{13}N_{15}N_{16})$$
(3.9)

where the first term is the usual MSSM contribution, the second and third terms are due to the singlet. The fourth and fifth terms are the triplet contribution coming from $\lambda_2 H_d T_0 H_u$ and $\lambda_4 STr(T_0 T_0)$ in the superpotential respectively. Numerical values of the components $S_{1j}(j = 1, ..., 4)$ as obtained from the benchmark point are, $S_{11} \sim 0.885$, $S_{12} \sim 0.463$, $S_{13} \sim 0.026$ and $S_{14} \sim -0.037$. In this model, we find that the spin-independent cross-section $\sigma_p \simeq 6.8 \times 10^{-10}$ pb (shown in Table.3.3), which is well below the upper bound presented by the latest XENON 100 results [55, 56] and can be accessible by the future XENON 1T experiment.

3.3 Di-photon Higgs decay rate in TSMSSM

In the SM, the diphoton decay of the Higgs boson is attributed through the Wboson loop and the contribution from the top-quark destructively interferes with the dominant W-boson contribution. The analytic expression for the diphoton partial width given as [154, 155]

$$\Gamma(h \to \gamma \gamma) = \frac{G_F \alpha^2 m_h^3}{128\sqrt{2}\pi^3} \left| A_1(\tau_W) + N_c Q_t^2 A_{1/2}(\tau_t) \right|^2 , \qquad (3.10)$$

where G_F is the Fermi constant, $N_c = 3$ is the number of color, $Q_t = +2/3$ is the top quark electric charge in units of |e|, and $\tau_i \equiv 4m_i^2/m_h^2$, i = t, W. The loop functions $A_1(\tau_W)$ and $A_{1/2}(\tau_t)$ for spin-1 (W boson) and spin-1/2 (top quark) particles are given in [156]. The numerical values of the loop functions for $m_h = 125$ GeV are,

$$A_1(\tau_W) \simeq -8.3 , A_{1/2}(\tau_t) \simeq 1.4$$

But in SUSY, we have additional contributions from the s-tops and charginos loops, which would significantly interfere with the SM contributions. Therefore, in general the branching width of Higgs decay to di-photon is formulated as [156],

$$\Gamma(h \to \gamma \gamma) = \frac{\alpha^2 m_h^3}{1024\pi^3} \left| \frac{g_{hVV}}{m_V^2} Q_V^2 A_1(\tau_V) + \frac{2g_{hf\bar{f}}}{m_f} N_{c,f} Q_f^2 A_{1/2}(\tau_f) + N_{c,S} Q_S^2 \frac{g_{hSS}}{m_S^2} A_0(\tau_S) \right|^2 ,$$
(3.11)

In the above the equation V, f, and S refer to generic spin-1, spin-1/2, and spin-0 particles, respectively. Q_V , Q_S and Q_f are the electric charges of the vectors, scalars and fermions in units of |e|, $N_{c,f}$ and $N_{c,S}$ are the number of fermion and scalar colors. $A_1(\tau_V)$, $A_{1/2}(\tau_f)$ and $A_0(\tau_S)$ are the loop functions for the vectors, fermions and scalars respectively.

In this model, the additional contribution to the diphoton Higgs decay width comes from the light chargino and the charged Higgs. Here, we take the assumption that the lightest charged Higgs (being dominantly triplet-like) only contribute to the decay width, since the other charged Higgs are much heavier. Now the term in the potential which gives rise to $hH^{\pm}H^{\pm}$ interaction is [145],

$$V_F \supset \lambda_2^2 v_u H_u^0 T_0^+ T_0^- \tag{3.12}$$

Therefore, the coupling $g_{hH^{\pm}H^{\pm}}$ becomes $\sim \lambda_2^2 v \sin \beta S_{11} C_{13} C_{14}$, where C_{ij} is the diagonalising matrix for the charged Higgs and $C_{13} \sim -0.669$, $C_{14} \sim -0.742$. The loop function for the scalar $A_0(\tau_s)$ is given by [156],

$$A_0(\tau_s) = -\tau_i^2 [\tau_s^{-1} - f(\tau_s^{-1})]]$$
(3.13)

where, $f(\tau_s) = arc \sin^2 \sqrt{\tau_s}$ for, $\tau_s > 1$.

Therefore, considering the main contributions due to charginos, charged triplet, W-boson and top quark t and in the limit $m_h^2 \ll 4m_{\tilde{\chi}_i^+}^2$, the diphoton Higgs decay


Figure 3.3: Top : Contours of $R_{\gamma\gamma}$ as a function of $\tan\beta$ and the triplet coupling λ_2 with $M_2 = 375$ GeV. Bottom : Contours of $R_{\gamma\gamma}$ as a function of $\tan\beta$ and M_2 with $\lambda_2 = 0.8$

rate with respect to the SM value becomes [142],

$$R_{\gamma\gamma} = \left| 1 + \frac{\frac{4}{3} \frac{\partial}{\partial \log v} \log \det \mathcal{M}_{ch}(v) + \frac{g_{hH\pm H\pm}}{m_{H_{1}}^{2}} A_{0}(\tau_{s})}{A_{1}(\tau_{W}) + \frac{4}{3} A_{1/2}(\tau_{t})} \right|^{2} , \qquad (3.14)$$

The numerator (first term) in Eq. (3.14) is given by

$$\frac{\partial}{\partial \log v} \log \det \mathcal{M}_{ch}(v) = -\frac{v^2 A}{2(M_2 \lambda_4 v_s + g_2^2 v_t^2) \mu_{eff} - \frac{1}{2} v^2 A},$$
(3.15)

where, $A = [\sin 2\beta(\lambda_2^2 M_2 + 2g_2^2\lambda_4 v_s) - 2\lambda_2 g_2^2 v_t]$ and the sign depends on the specific choices for the parameters. We are specifically interested in the region of parameter space where the numerator is negative (since the denominator is also negative), such that we obtain, $R_{\gamma\gamma} > 1$. We find that, the factor $g_{hH^{\pm}H^{\pm}}/m_{H_1^{\pm}}^2 \sim 0.0024$ and thus the contribution due to the extra charged triplet is treated to be negligible compared to the light chargino loops.

We see that for the set of benchmark points specified in Table.I, we obtain chargino masses in the range, $M_{\chi_i^{\pm}} \ni [131.92,299.38,422.24]$ GeV for tan $\beta = 1.8$. This choice of parameter gives, $R_{\gamma\gamma} \simeq 1.224$. In FIG.3.3 (left panel), we show the contours of $R_{\gamma\gamma}$ in the $(\tan\beta,\lambda_2)$ plane for $M_2 = 375$ GeV. We observe that 50% enhancement can be achieved with $\tan\beta \simeq 2$ but the triplet coupling λ_2 (≥ 1.1) then enters into nonperturbative regime. Right panel of FIG.3.3 shows the dependence of $R_{\gamma\gamma}$ on $\tan\beta$ and M_2 . Here, we note that lowering the value of M_2 increases the $R_{\gamma\gamma}$, but then we deviate from other phenomenological requirements.

3.4 Conclusion and Outlook

Recent analysis of the Fermi-LAT data shows existence of a monochromatic γ -ray line like features at $E_{\gamma} \sim 130$ GeV in the vicinity of the galactic center. A possible interpretation comes from DM annihilation into two photons, which requires the annihilation cross-section to be 1.27×10^{-27} cm³sec⁻¹ [118–121]. In this chapter, we have performed a scan over the parameter space of TSMSSM model and choose a specific set of benchmark points such that it satisfies all phenomenological requirements in order to obtain the required cross-section through the pair annihilation of 130 GeV neutralino DM via a psedoscalar Higgs triplet of mass $M_{A_T} \sim 2m_{DM}$ near resonance and light chargino loops. The width of the pseudoscalar triplet being small helps in boosting the Breit-Weigner cross-section, $\langle \sigma v \rangle_{\gamma\gamma}$. Besides, this model also predicts a second γ -ray peak at 114 GeV from the annihilation $\chi\chi \to \gamma Z$, and the cross section is approximately 0.75 times that of $\langle \sigma v \rangle_{\gamma\gamma}$, which is below the upper limit reported by Fermi LAT. The dark matter candidate being a mixture of bino-higgsino, leads to a correct relic abundance of 0.109, consistent with the PLANCK and WMAP-9 year data. The spin-independent scattering cross-section with nucleons is 0.68×10^{-9} pb, which is well below the latest XENON100 exclusion limits.

Although latest results from CMS seem to favour a SM-like Higgs boson, but on the other hand ATLAS still shows a significant excess in diphoton decay width compared to SM as, $R_{\gamma\gamma} = 1.65^{+0.34}_{-0.30}$ for $m_h = 126$ GeV. Our model predicts a similar enhancement in the diphoton decay rate as, $R\gamma\gamma \sim 1.224$, which is contributed dominantly through the light chargino loops, since the contribution from the extra charged triplet is negligible. Such a prediction opens the possibility of this model being tested in future LHC runs.

Chapter 4

Constraining Minimal $U(1)_{B-L}$ Model from Dark matter observations

4.1 Overview

Many extensions of the SM were proposed with a motivation to introduce a suitable DM candidate. Among the plethora of candidates, the weakly interacting massive particles (WIMP) are the popular choice (for review see [21, 22, 29]). A simplest extension of the SM with a real or complex gauge singlet scalar field [42-47] (see Chap.1) has been extensively studied. The scalar turns out to be an appropriate DM candidate, which interacts only with the SM Higgs boson. Another possibility includes a renormalizable extension of the SM with a gauge singlet Dirac fermion (ψ) along with a gauge singlet scalar (S) [48– 53] (see Chap.1), known as Singlet Fermionic Dark Matter (SFDM) model. In SFDM, the singlet scalar interact with the SM Higgs boson whereas ψ becomes the viable DM candidate, which interacts to the SM particles via S only. On the other hand, neutrino mass generation can be linked with DM mass through the radiative seesaw mechanism [157–159], and the Ma-model [160]. Among other possibilities, the minimal gauge extension of the SM with $U(1)_{B-L}$, and a discrete symmetry (Z_2 -parity) has been studied by several authors [157–159, 161–163] in the context of DM.

In this chapter, we study the minimal $U(1)_{B-L}$ extension of the SM [164– 166], with an additional Z_2 -symmetry imposed on the model [161]. Here, only one of the right-handed (RH) neutrinos being odd under Z_2 -parity, serves as an excellent DM candidate. We obtain effectively a Higgs-portal DM which can annihilate into the SM particles (dominantly into W^+W^- and ZZ) and gives correct relic abundance [4, 5] near resonances where DM mass is almost half of the scalar boson masses. Our primary motivation is to restrict the choice of parameter space of this model, based on various recent experimental results of dark matter like relic abundance, limits on spin-independent scattering crosssection etc, which has not been considered in earlier studies. We emphasize that the heavy scalar decay width depends strongly on the scalar mixing angle and hence plays a significant role in determining the relic density. Demanding correct relic abundance we constrain the parameter space of the scalar mixing angle and heavy scalar boson mass. We found that the spin-independent elastic scattering cross-section off nucleon is maximum at a particular value of scalar mixing angle and lies below the XENON100 [55, 56] and the latest LUX [57] exclusion limits. However the future XENON1T [54] experiment can further restrict the heavy scalar mass. Using the constraints on scalar mixing angle and heavy scalar mass, we have also calculated the annihilation cross-section into two photon final state $\langle \sigma v \rangle_{\gamma\gamma}$ and finally compare with the upper bound on $\langle \sigma v \rangle_{\gamma\gamma}$ by Fermi-LAT [167] for different DM profiles. We observe that the resultant $\langle \sigma v \rangle_{\gamma\gamma}$ coincide with the Fermi-LAT data in the region where DM mass is almost half of the light scalar boson mass, otherwise it is well below the Fermi-LAT bound. Apart from DM phenomenology, neutrino mass can be generated in this model via Type-I seesaw mechanism. Here the lightest neutrino remains massless (because of odd- Z_2 parity of one of the RH-neutrinos), which is consistent with the observed oscillation data.

The chapter is organized as follows: The next section contains a brief description of the model; we discuss the observational constraint from dark matter in Section 4.3 with an estimation of the relic density in Section 4.3.1, direct detection of the DM in Section 4.3.2 and a detail calculation for annihilation into two

Particle	Q	u_R	d_R	L	e_R	Φ	S	$N_{R^{1,2}}$	N_{R^3}
$SU(2)_L$	2	1	1	2	1	2	1	1	1
$U(1)_Y$	1/6	2/3	-1/3	-1	-1	1	0	0	0
$U(1)_{B-L}$	1/3	1/3	1/3	-1	-1	0	2	-1	-1
\mathbb{Z}_2	+	+	+	+	+	+	+	+	-

Table 4.1: Particle content of minimal $U(1)_{B-L}$ model

photon final state can be found in Section. 4.3.3; finally we present the results and analysis in Section 4.4 and summarize our results and conclude in the last section.

A detail calculation of the total decay width of the heavy scalar boson has been shown in appendix A. Appendix B shows the estimation of w(s) required for the calculation of relic abundance. Appendix C contains the loop functions necessary for calculating the cross-sections $\langle \sigma v \rangle_{\gamma\gamma}$.

4.2 Minimal Gauged $U(1)_{B-L}$ Model

In this work, we adopt the minimal $U(1)_{B-L}$ extension of the SM [164–166]. Along with the SM particles, this model contains: a SM singlet S with B - L charge +2, three right-handed neutrinos $N_R^i(i = 1, 2, 3)$ having B - L charge -1. As this $U(1)_{B-L}$ symmetry is gauged, an extra gauge boson Z' is associated as a signature of the extended symmetry. Once the B - L symmetry is broken spontaneously through the vacuum expectation value (vev) of S, this Z' becomes massive. Here, we also impose a Z_2 discrete symmetry. We assign Z_2 charge +1(or even) for all the particles except N_R^3 [161, 168]. This ensures the stability of N_R^3 which qualified as a viable DM candidate. The assignment of B-L charge in this model eliminates the triangular B - L gauge anomalies and ensures the gauge invariance of the theory. The particle content of this model is shown in Table 4.1. Scalar Lagrangian of this model can be written as,

$$\mathcal{L}_{s} = (D^{\mu}\Phi)^{\dagger} D_{\mu}\Phi + (D^{\mu}S)^{\dagger} D_{\mu}S - V(\Phi, S), \qquad (4.1)$$

where the potential term is,

$$V(\Phi, S) = m^2 \Phi^{\dagger} \Phi + \mu^2 |S|^2 + \lambda_1 (\Phi^{\dagger} \Phi)^2 + \lambda_2 |S|^4 + \lambda_3 \Phi^{\dagger} \Phi |S|^2, \quad (4.2)$$

with Φ and S as the Higgs doublet and singlet fields, respectively. After spontaneous symmetry breaking (SSB) the two scalar fields can be written as,

$$\Phi = \begin{pmatrix} 0\\ \frac{v+\phi}{\sqrt{2}} \end{pmatrix}, \qquad S = \frac{v_{B-L} + \phi'}{\sqrt{2}}, \qquad (4.3)$$

with v and v_{B-L} real and positive. Minimization of eq. (4.2) gives

$$m^{2} + 2\lambda_{1}v^{2} + \lambda_{3}vv_{B-L}^{2} = 0,$$

$$\mu^{2} + 4\lambda_{2}v_{B-L}^{2} + \lambda_{3}v^{2}v_{B-L} = 0.$$
(4.4)

To compute the scalar masses, we must expand the potential in eq. (4.2) around the minima in eq. (4.3). Using the minimization conditions, we have the following scalar mass matrix :

$$\mathcal{M} = \begin{pmatrix} \lambda_1 v^2 & \frac{\lambda_3 v_{B-L} v}{2} \\ \frac{\lambda_3 v_{B-L} v}{2} & \lambda_2 v_{B-L}^2 \end{pmatrix} = \begin{pmatrix} \mathcal{M}_{11} & \mathcal{M}_{12} \\ \mathcal{M}_{21} & \mathcal{M}_{22} \end{pmatrix}.$$
 (4.5)

The expressions for the scalar mass eigenvalues $(m_H > m_h)$ are:

$$m_{H,h}^2 = \frac{1}{2} \bigg[\mathcal{M}_{11} + \mathcal{M}_{22} \pm \sqrt{(\mathcal{M}_{11} - \mathcal{M}_{22})^2 + 4\mathcal{M}_{12}^2} \bigg].$$
(4.6)

The mass eigenstates are linear combinations of ϕ and ϕ' , and written as

$$\begin{pmatrix} h \\ H \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} \phi \\ \phi' \end{pmatrix}, \quad (4.7)$$

where, h is the SM-like Higgs boson. The scalar mixing angle, α can be expressed as:

$$\tan(2\alpha) = \frac{2\mathcal{M}_{12}}{\mathcal{M}_{11} - \mathcal{M}_{22}} = \frac{\lambda_3 v_{B-L} v}{\lambda_1 v^2 - \lambda_2 v_{B-L}^2}.$$
 (4.8)

Now we can calculate the quartic coupling constants by using eqs. (4.6, 4.7 and 4.8).

$$\lambda_{1} = \frac{m_{H}^{2}}{4v^{2}}(1 - \cos 2\alpha) + \frac{m_{h}^{2}}{4v^{2}}(1 + \cos 2\alpha),$$

$$\lambda_{2} = \frac{m_{h}^{2}}{4v_{B-L}^{2}}(1 - \cos 2\alpha) + \frac{m_{H}^{2}}{4v_{B-L}^{2}}(1 + \cos 2\alpha),$$

$$\lambda_{3} = \sin 2\alpha \left(\frac{m_{H}^{2} - m_{h}^{2}}{2 v v_{B-L}}\right).$$
(4.9)

In the presence of an extra $U(1)_{B-L}$ gauge theory the SM gauge kinetic terms is modified by

$$\mathcal{L}_{B-L}^{K.E} = -\frac{1}{4} F'^{\mu\nu} F'_{\mu\nu} , \qquad (4.10)$$

where,

$$F'_{\mu\nu} = \partial_{\mu}B'_{\nu} - \partial_{\nu}B'_{\mu}. \qquad (4.11)$$

The general covariant derivative in this model reads as

$$D_{\mu} \equiv \partial_{\mu} + ig_{S}T^{\alpha}G_{\mu}^{\ \alpha} + igT^{a}W_{\mu}^{\ a} + ig_{1}YB_{\mu} + i(\tilde{g}Y + g_{B-L}Y_{B-L})B_{\mu}'.$$
(4.12)

Here, we consider only the 'pure' B - L model, that is defined by the condition $\tilde{g} = 0$ at Electro-Weak (EW) scale. This implies zero mixing at tree level between Z' and Z bosons.

The relevant Yukawa coupling to generate neutrino masses is given by,

$$\mathcal{L}_{int} = \sum_{\beta=1}^{3} \sum_{j=1}^{2} y_{\beta}^{j} \overline{l_{\beta}} \tilde{\Phi} N_{j} - \sum_{i=1}^{3} \frac{y_{n_{i}}}{2} \overline{N_{R}^{i}} S N_{R}^{i}$$
(4.13)

where, $\tilde{\Phi} = -i\tau_2 \Phi^*$.

The neutrino mass can be generated in this model via Type-I seesaw mechanism, where the mass matrices for light and heavy neutrino are given as,

$$m_{\nu_L} \simeq m_D^T m_M^{-1} m_D,$$
 (4.14)

$$m_{\nu_H} \simeq m_M \tag{4.15}$$

where, $m_D = (y_{\beta}^j/\sqrt{2})v$, (j = 1, 2) and $m_{M_i} = -(y_{n_i}/\sqrt{2})v_{B-L}$, (i = 1, 2, 3).

Because of Z_2 -parity, N_R^3 has no Yukawa coupling with the left-handed lepton doublet, therefore the lightest neutrino remains massless. The masses of N_R^1 and N_R^2 are considered to be heavier than that of N_R^3 .

4.3 Dark Matter Observations

4.3.1 Relic Density

In the early universe when the temperature was high enough, the DM particles were in thermal equilibrium with the rest of the cosmic plasma and its number density had fallen off exponentially with temperature. But as temperature dropped down below the DM mass, the annihilation rate decreased and became smaller than the Hubble expansion rate. Then the DM species was decoupled from the cosmic plasma and number density experienced a "freeze-out" - hence we observe a significant relic abundance of DM today.

The relic abundance of DM can be formulated as [30],

$$\Omega_{CDM}h^2 = 1.1 \times 10^9 \frac{x_f}{\sqrt{g^* m_{Pl} \langle \sigma v \rangle_{ann}}} \text{GeV}^{-1} , \qquad (4.16)$$

where $x_f = m_{N_R^3}/T_D$ with T_D as decoupling temperature. m_{Pl} is Planck mass = 1.22×10^{19} GeV, and, g^* is effective number of relativistic degrees of freedom (we use, $g^* = 100$ and $x_f = (1/20)$). $\langle \sigma v \rangle_{ann}$ is the thermal averaged value of DM annihilation cross-section times relative velocity. DM interacts with the SM particles via Z'-boson and h, H. But, Z'-boson being heavy ($m_{Z'} \ge 2.33$ TeV [169]), the annihilation of DM into the SM particles takes place via h and Honly. Thus, effectively we obtain a Higgs-portal DM model.

 $\langle \sigma v \rangle_{ann}$ can be obtained using the well known formula [170],

$$\langle \sigma v \rangle_{ann} = \frac{1}{m_{N_R^3}^2} \left\{ w(s) - \frac{3}{2} \left(2w(s) - 4m_{N_R^3}^2 w'(s) \right) \frac{1}{x_f} \right\} \Big|_{s = \left(2m_{N_R^3}^2 \right)^2} , \qquad (4.17)$$

where prime denotes differentiation with respect to s (\sqrt{s} is the center of mass energy). Here, the function w(s) (detail calculation in appendix B) depends on amplitude of different annihilation processes,

$$N_R^3 N_R^3 \longrightarrow b\bar{b}, \ \tau^+ \tau^-, \ W^+ W^-, \ ZZ, \ hh.$$
 (4.18)

4.3.2 Spin-independent scattering cross-section

The effective Lagrangian describing the elastic scattering of the DM off a nucleon is given by,

$$L_{eff} = f_p \bar{N}_R^3 N_R^3 \bar{p}p + f_n \bar{N}_R^3 N_R^3 \bar{n}n , \qquad (4.19)$$

where, $f_{p,n}$ is the hadronic matrix element, given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q}.$$
 (4.20)

The f-values are given as in [107]

$$f_{Tu}^{(p)} = 0.020 \pm 0.004, \quad f_{Td}^{(p)} = 0.026 \pm 0.005, \quad f_{Ts}^{(p)} = 0.118 \pm 0.062 ,$$

$$f_{Tu}^{(n)} = 0.014 \pm 0.003, \quad f_{Td}^{(n)} = 0.036 \pm 0.008, \quad f_{Ts}^{(n)} = 0.118 \pm 0.062 ,$$

and $f_{TG}^{(p,n)}$ is related to these values by

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}.$$
(4.21)

Here, a_q is the effective coupling constant between the DM and the quark. We obtain the scattering cross-section (spin-independent) for the dark matter off a proton or neutron as,

$$\sigma_{p,n}^{SI} = \frac{4m_r^2}{\pi} f_{p,n}^2 \tag{4.22}$$

where, m_r is the reduced mass defined as, $1/m_r = 1/m_{N_R^3} + 1/m_{p,n}$.

An approximate form of a_q/m_q can be recast in the following form :

$$\frac{a_q}{m_q} = \frac{y_{n_3}}{v\sqrt{2}} \left[\frac{1}{m_h^2} - \frac{1}{m_H^2} \right] \sin\alpha \cos\alpha , \qquad (4.23)$$

where, $y_{n_3} = \sqrt{2} m_{N_R^3} / v_{B-L}$ is the Yukawa coupling as specified in the second term of eq. (4.13).

4.3.3 Annihilation cross-section into two photons

The RH-neutrino dark matter N_R^3 can also annihilate into two photon final state mediated by scalar bosons (*h* and *H*) through loop suppressed processes. Here, we consider mostly dominant contributions from top-quark and W-boson loops to this process [53].

The thermal averaging of the annihilation cross-section $\sigma v_{\gamma\gamma}$ can be obtained using [170]

$$\langle \sigma v \rangle_{\gamma\gamma} = \frac{1}{m_{N_R^3}^2} \left\{ w(s)_{\gamma\gamma} - \frac{3}{2} \left(2w(s)_{\gamma\gamma} - 4m_{N_R^3}^2 w'(s)_{\gamma\gamma} \right) \frac{1}{x_f} \right\} \Big|_{s = \left(2m_{N_R^3} \right)^2} .$$
(4.24)

The function $w(s)_{\gamma\gamma}$ for massless final product is defined as,

$$w(s)_{\gamma\gamma} = \frac{1}{32\pi} \sum_{spins} |\mathcal{M}_{N_R^3 N_R^3 \to \gamma\gamma}|^2.$$
(4.25)

Taking into account contributions via h and H bosons we obtain,

$$\sum_{spins} |\mathcal{M}_{N_{R}^{3}N_{R}^{3} \to \gamma\gamma}|^{2} = y_{n_{3}}^{2} (s - 4m_{N_{R}^{3}}^{2}) \left\{ \frac{|\mathcal{M}_{h \to \gamma\gamma}|^{2} \sin^{2} \alpha}{(m_{h}^{2} - s)^{2} + m_{h}^{2}\Gamma_{h}^{2}} + \frac{|\mathcal{M}_{H \to \gamma\gamma}|^{2} \cos^{2} \alpha}{(m_{h}^{2} - s)^{2} + m_{h}^{2}\Gamma_{h}^{2}} + \frac{|\mathcal{M}_{h \to \gamma\gamma}||\mathcal{M}_{H \to \gamma\gamma}|\sin\alpha\cos\alpha\{(m_{h}^{2} - s)(m_{H}^{2} - s) + m_{h}m_{H}\Gamma_{h}\Gamma_{H}\}}{((m_{h}^{2} - s)^{2} + m_{h}^{2}\Gamma_{h}^{2})((m_{H}^{2} - s)^{2} + m_{H}^{2}\Gamma_{H}^{2})} \right\}.$$

$$(4.26)$$

where, $\mathcal{M}_{h(H)\to\gamma\gamma}$ is the amplitude for the decay of h(H) into two photons, which reads as [171, 172]

$$\mathcal{M}_{h(H)\to\gamma\gamma} = \frac{g_2 \ \alpha_{\rm em} \ m_{h,H}^2}{8\pi m_W} \Big[3\left(\frac{2}{3}\right)^2 F_t(\tau_t) + F_W(\tau_W) \Big] \cos\alpha(\sin\alpha) , \qquad (4.27)$$

where, $\tau_i = 4m_i^2/m_{h,H}^2$ (i = W, t) and $F_{W,t}(\tau_{W,t})$ are the loop functions for Wboson and top-quark respectively (see appendix C for detail calculation). $\alpha_{\rm em}$ is the electromagnetic fine structure constant at the EW scale, $\alpha_{\rm em}(m_Z) \sim 1/127$. SU(2) gauge coupling is denoted as g_2 , whereas, m_W is the W-boson mass.

4.4 **Results and Analysis**

In this model, the right-handed neutrino N_R^3 turns out to be a viable dark matter candidate as an artifact of the Z_2 charge assignment. We choose a specific set of benchmark values for (mass (m_h) and decay width (Γ_h) of SM-like Higgs boson, *vev* of singlet scalar S and $U(1)_{B-L}$ gauge coupling) our calculation, shown in

m_h	Γ_h	$v_{\scriptscriptstyle B-L}$	$g_{\scriptscriptstyle B-L}$
$125~{\rm GeV}$	$4.7 \times 10^{-3} \text{ GeV}$	$7 { m TeV}$	0.1

 Table 4.2: Choice of Parameters



Figure 4.1: Plot of relic abundance as a function of DM mass for $m_H = 500 \text{ GeV}$ with specific choices of scalar mixing angle $\cos \alpha = 0.935$ (blue-dashed), 0.45 (redsolid). The straight line shows the WMAP9 value, $\Omega_{CDM} h^2 = 0.1148 \pm 0.0019$.

Table.4.2, based on present experimental constraints [169]. However, the mass of the heavy scalar and the scalar mixing angle are not fixed.

In Figure. 4.1 the relic density is plotted against DM mass for two specific choices (to be explained later in this section) of scalar mixing angles $\cos \alpha = 0.935$, 0.45 with $m_H = 500$ GeV. The straight line shows the latest 9-year WMAP data i.e, $\Omega_{CDM}h^2 = 0.1148 \pm 0.0019$ [4] (whereas latest PLANCK result is, $\Omega_{CDM}h^2 = 0.1199 \pm 0.0027$ at 68% CL [5]). The resultant relic abundance is found to be consistent with the reported value of WMAP-9 and PLANCK experiment only near resonance when, $m_{N_R^3} \sim (1/2) m_{h,H}^{-1}$. The reason for

¹In principle, Z' resonance can also provide the correct relic abundance, but in that case



Figure 4.2: Variation of w'(s) near resonances : (a) $m_{N_R^3} = m_h/2$ and (b) $m_{N_R^3} = m_H/2$, with $m_h = 125$ GeV and $m_H = 500$ GeV, respectively.

the over abundance of DM except at the resonance can be understood in the following way : The annihilation cross-section of DM, being proportional to $y_{n_3}^2$ (where, $y_{n_3} = (\sqrt{2}m_{N_R^3})/v_{B-L}$), is heavily suppressed due to large value of v_{B-L} . Figure. 4.1 also exhibits a strong dependence on the mixing angle near the second

the DM mass will be $\mathcal{O}(\text{TeV})$ (i.e $m_{N_R^3} \sim (1/2) m_{Z'}$), if we consider the current experimental bound on Z' mass [169].



Figure 4.3: Yellow region (in the middle) shows the allowed range of $\cos \alpha$ and m_H consistent with correct relic abundance as reported by WMAP9. The above-pink (below-white) region is disallowed due to under-abundance (over-abundance) of dark matter.

resonance (i.e, $m_{N_R^3} \sim (1/2) m_H$). Since, the criterion for correct relic abundance is satisfied near scalar resonances, we have studied the contribution of different annihilation channels to the total annihilation cross-section in that region. We have plotted in Figure. 4.2 the variation of w'(s) ($\langle \sigma v \rangle_{ann}$ depends on w'(s) as shown in eq. (4.17)) near resonances $m_{N_R^3} = m_{h,H}/2$ for different annihilation channels like $b\bar{b}$, $\tau^+\tau^-$, W^+W^- , ZZ, hh. We observe that the dominant contribution to the total annihilation cross-section comes from the W^+W^- , ZZ (also final state hh dominance observed in Figure. 4.2(b)) final states, which is expected because of large SU(2) gauge coupling. In case of Figure. 4.2(a) a sharp (narrow) resonance peak is observed, whereas figure. 4.2(b) has a broad resonance due to larger decay width (Γ_H) of the heavy scalar, which also depends on scalar mixing angle (see appendix A).

Relic abundance near the second resonance depends on the following model parameters (unknown) : scalar mixing angle (α), heavy scalar mass (m_H) and



Figure 4.4: Variation of σ_p^{SI} with $m_{N_R^3}$ for $m_H = 300$ GeV (green-dashed) and 900 GeV (black-solid) with $\cos \alpha = 0.707$. The blue and violet curves show the bound from XENON100 [55, 56] and LUX [57] data respectively. Red curve shows the projected limits for XENON1T [54].

decay width (Γ_H). But, these are not independent as Γ_H can be derived using cos α and m_H . For large mixing angle, the total decay width of heavy scalar is large and hence the annihilation cross-section $\langle \sigma v \rangle_{ann}$ is less compared to that with minimal mixing scenario. This behavior is observed in Figure. 4.1, where $\Omega_{CDM}h^2$ is large for smaller value of cos α (at $m_{N_R^3} \sim (1/2) m_H$) and vice-versa. We therefore perform a scan over the entire parameter range of m_H (300-1000 GeV) and cos α [168] to find the allowed region consistent with the 9-year WMAP data ($\Omega_{CDM}h^2 = 0.1148 \pm 0.0019$)[4]. In Figure. 4.3, the yellow region shows the allowed (by correct relic abundance) range of cos α for different values of m_H , whereas the pink region is forbidden because the annihilation cross-section is enhanced for smaller mixing angle (smaller decay width Γ_H) leading to underabundance of dark matter. On the other hand, the white region is disallowed because of over-abundance.

According to eq. (4.22), as shown in section 4.3.2, it is evident that, $\sigma_{p,n}^{SI} \propto (\sin 2\alpha)^2 f(m_H)$, which is maximum at $\alpha = \pi/4$ (or $\cos \alpha = 0.707$) irrespective



Figure 4.5: Annihilation cross-section into two photon final state vs. dark matter mass with two specific choices : $\cos \alpha = 0.935, m_H = 500$ GeV (bluesolid) and $\cos \alpha = 0.885, m_H = 390$ GeV (purple-dashed) respectively. The upper-most two curves show the Fermi-LAT upper bound on $\langle \sigma v \rangle_{\gamma\gamma}$ [167] for NFW (solid-red) and Einasto (dashed-black) profile.

of the choice of m_H . Therefore, the maximum value of $\sigma_{p,n}^{SI}$ increases as m_H is increased, which can be understood from eqs. (4.22, 4.23). Figure. 4.4 shows [168] the maximum value of spin-independent scattering cross-section (i.e, with $\cos \alpha = 0.707$) of the DM off proton (σ_p^{SI}) for $m_H = 300$ GeV (green-dashed) and 900 GeV (black-solid), whereas the blue and violet curves show the XENON100 (2012) [55, 56] and the latest LUX (at 95% C.L.) [57] exclusion plots, respectively. The red-curve shows the projected limits on σ_p^{SI} for XENON1T experiment [54]. We observe that the value of the resultant cross-section with two different values of m_H for the entire range 6 GeV $\leq m_{N_R^3} \leq 500$ GeV lies much below the XENON100 and latest LUX exclusion limits. But, as the value of m_H is increased, the spin-independent cross-section becomes larger at higher values of DM mass and approaches the limits as reported by LUX and XENON100. As shown in Figure. 4.4, in future XENON1T data might severely restrict the choice of allowed m_H . As described in section 4.3.3, figure. 4.5 shows the maximum annihilation cross-section into two photon final state as a function of dark matter mass with different values of $\cos \alpha$ and m_H . Here, we have chosen the maximum allowed value of $\cos \alpha$ corresponding to particular value of m_H as derived in Section. 4.4 (see Figure. 4.3). The blue(pink-dashed) curve shows the resultant $\langle \sigma v \rangle_{\gamma\gamma}$ for $\cos \alpha = 0.935(0.885)$ and $m_H = 500(390)$ GeV. It also shows a comparison [168] with the Fermi-LAT upper bound on $\langle \sigma v \rangle_{\gamma\gamma}$ for Navarro-Frenk-White (NFW) (solid-red) and Einasto (dashed-black) profile [167]. We observe a clear coincidence between theoretical plots and Fermi-LAT data near resonance point where $m_{N_R^3} \sim (1/2) m_h$. A second peak is observed in the pink-curve due to a second resonance at $m_{N_R^3} \sim (1/2) m_H$ (i.e. at 195 GeV), but the maximum $\langle \sigma v \rangle_{\gamma\gamma}$ is found to be much below the exclusion limit of Fermi-LAT data.

Comment on 130 GeV monochromatic γ -ray line : Last year, the analysis of Fermi-LAT data [68] had revealed a hint of a monochromatic gamma ray features [118, 119, 121] with $E_{\gamma} \simeq 130$ GeV coming from the vicinity of Galactic Center. One of the possible explanations of this phenomena could arise from the annihilation of DM with mass $129.8 \pm 2.4^{+7}_{-13}$ GeV and annihilation cross-section $\langle \sigma v \rangle_{\gamma\gamma} = (1.27 \pm 0.32^{+0.18}_{-0.28}) \times 10^{-27} cm^3 sec^{-1}$. It is possible to explain this monochromatic photon line in this model with a resonant heavy scalar near 260 GeV and achieve the desired cross-section. But, since the DM dominantly annihilates into W^+W^- , ZZ final states ($\langle \sigma v \rangle_{\gamma\gamma}$ is also suppressed as $\mathcal{O}(\alpha_{em}^2(M_Z))$, the continuum photon spectra supersaturate the monochromatic line-like feature.

4.5 Conclusion and Outlook

In this chapter, we have studied a minimal $U(1)_{B-L}$ extended SM, where the third generation RH-neutrino becomes the plausible DM candidate by the virtue of an additional Z_2 -symmetry. The DM considered in this model is effectively Higgs-portal and annihilates dominantly into gauge boson (W^+W^-, ZZ) final states. We derive an important constraint on the allowed parameter space of the scalar mixing angle and heavy scalar mass in order to obtain correct relic abundance. Besides this, the relic abundance is found to be consistent with the recent WMAP9 and PLANCK data only near scalar resonances, i.e., $m_{N_{P}^{3}} = (1/2) m_{h,H}$.

WMAP9 and PLANCK data only near scalar resonances, i.e, $m_{N_R^3} = (1/2) m_{h,H}$. In future, PLANCK data can further restrict the choice of parameter space. The total annihilation cross-section is enhanced due to scalar resonance, otherwise it will be suppressed due to heavy Z'. The spin-independent elastic scattering cross-section of DM off a nucleon is maximum for $\cos \alpha = 0.707$, and hence maximum σ_p^{SI} depends on the value of heavy scalar mass. We observe that, σ_p^{SI} is well below the XENON100 and LUX exclusion limits for DM mass ranging from 5-500 GeV. But, future direct detection experiments like XENON1T can put stringent constraint on the choice of m_H . The annihilation cross-section of dark matter into $\gamma\gamma$ mediated by h and H bosons is compared with that of Fermi-LAT upper bound. We find an agreement between the theoretical plot and the Fermi-LAT data near scalar resonance where, $m_{N_R^3} = (1/2) m_h$. Although the required $\langle \sigma v_{\gamma\gamma} \rangle$ for explaining 130 GeV Fermi-line can be obtained in this model, but the gamma-ray continuum spectra produced due to W^+W^- , ZZ final state supersaturate this monochromatic line feature. In addition, this model can successfully account for the neutrino masses generated via Type-I seesaw mechanism. In future, more precise determination of relic abundance and scattering cross-section can be used for obtaining stronger bounds on the allowed parameter space of this kind of model.

Chapter 5

Future Directions and Discussions

In this present thesis, we have explored both supersymmetric and non-supersymmetric extension of SM, in order to accommodate a suitable DM candidate. The main issue of DM is to reconcile two different cross-sections - total annihilation crosssection required for relic abundance is $\sim 10^{-26} \mathrm{cm}^3 \mathrm{sec}^{-1}$ and exclusion limit on spin-independent scattering cross-section suggests $\sigma_p^{SI} \sim 10^{-45} {\rm cm}^2.$ A successful DM model must fulfill these requirements. However, in models with inelastic DM, one can evade the bound on scattering cross-section of DM off nucleon. We have introduced a new supersymmetric model, the Triplet-Singlet extension of MSSM, where the components of the additional singlet and triplet superfields play an important role in both Higgs sector and fermionic sector. In this model, Higgs mass can reach up to 122 GeV at the tree-level and does not require large radiative corrections, which results in less fine-tuning compared to other SUSY models. The neutralino-LSP turns out to be a viable DM candidate. We have intensively studied the DM phenomenology of this model and explained the monochromatic γ -ray line-like feature. We have also addressed the reported enhancement in the di-photon Higgs decay rate. Dark matter in triplet-SUSY models have not been much explored in the past. In future, we wish to implement this model in the numerical codes like SARAH [173] and micrOMEGAs [149] - which would enable us to study other DM and collider related phenomenology.

Apart from the Supersymmetric extension, we have adopted the minimal $U(1)_{B-L}$ extension of SM, with a RH-neutrino DM. We constrain the param-

eter space of this model subject to observational constraints on DM, like relic abundance, direct detection cross-section etc. The addition of a RH-neutrino as dark matter can also serve as a basis for neutrino mass models. One or more RH-neutrino with Majorana masses of ~1 TeV can be used in Seesaw models [174–176] where the leading order neutrino mass term, $M_D^T M_R^{-1} M_D = 0$ by symmetry and the next leading order term for neutrino mass eigenstates gives the observed light neutrino masses with a RH-neutrino (~ 1 TeV) as dark matter.

Gamma-ray emission from the galactic center (GC) and the inner galaxy regions as found in the Fermi-LAT data has gained a lot of attention from the perspective of dark matter (DM) searches. Past studies [177–183] have pointed out a spatially extended excess of $\sim 1-3$ GeV gamma rays from the regions surrounding the galactic center, the morphology and spectrum of which is best fitted with that predicted from the annihilations of a 31-40 GeV WIMP (weakly interacting massive particle) dark matter (DM) candidate annihilating mostly to b-quarks (or a $\sim 7 - 10$ GeV WIMP annihilating significantly to τ -leptons). Gamma rays from the galactic center is specially interesting because the region is predicted to contain very high densities of dark matter. Alternative explanations such as gamma-ray excess originating from thousands of unresolved millisecond pulsars have been disfavored since the signal extends well beyond the boundaries of the central stellar cluster. A more recent scrutiny of the morphology and spectrum of the anomalous gamma-ray emission in order to identify the origin has confirmed that the signal is very well fitted by a 31-40 GeV dark matter particle annihilating to $b\bar{b}$ with an annihilation cross section of $\sigma v = (1.4 - 1.4)$ 2.0) × 10⁻²⁶ cm³ sec⁻¹ (normalized to a local dark matter density of 0.3 GeV cm⁻³) [184], which is accidentally close to the weak cross-section for producing correct relic abundance. Already a handful of particle physics models of dark matter [185–204] have been proposed to explain the reported gamma-ray excess. In this context, we can study the Higgs-portal DM models which are well-suited for explaining these phenomena and constrain the parameter space of such models as a consequence [205].

In future, we would like to address some of the recent observations in the

context of DM - (i) positron excess seen by AMS-02, PAMELA [66, 67] and (ii) x-ray line emission in the energy range of a few keVs (for example see Ref. [206–209]). Now, the annihilation cross-section needed to explain positron excess, as seen in PAMELA, AMS-02 experiments, is around $\sim 10^{-24} \text{cm}^3 \text{sec}^{-1}$, which is a few order of magnitude larger than that required for obtaining correct relic abundance. Certainly, we require some kind of boost factor to accommodate this phenomena. Leptophilic DM models which can also evade direct detection bound, are suitable for explaining positron excess without any anti-proton excess.

The ratio of observed baryonic density to CDM density $\Omega_b/\Omega_{CDM} \sim 1/5$ which implies that with equal number densities $n_b \simeq n_{CDM}$, the dark matter mass would be of the order of $5m_p \sim 5$ GeV. The number densities of baryon and CDM (which are unrelated in the thermal relic density mechanism) can be related if they have a common origin i.e., they are produced in the same process. Such scenarios are called cogenesis [210–212] and in these cases the dark matter annihilation cross-section is not restricted to the relic density constraints. With DM mass of the order of \sim 5 - 10 GeV one can explain also the galactic centre gamma-ray excess. Also a common origin of DM of baryon would imply that there is an excess of particles over anti-particles in both DM and baryonic sectors and these scenarios are called Asymmetric dark matter (aDM) (for review see Ref. [213, 214]). If there is an asymmetry in the dark sector, as soon as annihilations have wiped out the density of antiparticles, the number density of particles remains frozen for lack of targets, and is entirely controlled by the primordial asymmetry rather than by the thermal freeze-out. In future, we would also like to explore the phenomenology associated to oscillations between DM/anti-DM [215].

Appendix A

Decay width of heavy scalar boson

In this model we have two Higgs mass eigenstates (h, H) which are admixture of the gauge eigenstates with the mixing angle α . The SM gauge eigenstate (ϕ) can be written as

$$\phi = \cos \alpha \ h + \sin \alpha \ H$$

So the coupling of h(H) with the SM particles will be multiplied by $\cos\alpha(\sin\alpha)$.

Decay of heavy scalar into fermion–antifermion (SM) pair

$$\Gamma(H \to f\bar{f}) = N_c \frac{g^2 m_f^2 m_H}{32 \pi m_W^2} \left\{ 1 - \frac{4m_f^2}{m_H^2} \right\}^{3/2} (\sin\alpha)^2$$
(A.1)

where N_c is the color factor, 1 for leptons and 3 for quarks.

Decay of heavy scalar into W boson pair

$$\Gamma(H \to W^+ W^-) = \frac{g^2 m_H^3}{64 \pi m_W^2} \sqrt{1 - \frac{4m_W^2}{m_H^2}} \left[1 - \frac{4m_W^2}{m_H^2} + \frac{3}{4} \left(\frac{4m_W^2}{m_H^2} \right)^2 \right] (\sin \alpha)^2$$
(A.2)

Decay of heavy scalar into Z boson pair

$$\Gamma(H \to ZZ) = \frac{g^2 m_H^3}{128 \pi m_W^2} \sqrt{1 - \frac{4m_Z^2}{m_H^2}} \left[1 - \frac{4m_Z^2}{m_H^2} + \frac{3}{4} \left(\frac{4m_Z^2}{m_H^2} \right)^2 \right] (\sin \alpha)^2$$
(A.3)

Decay of heavy scalar into RH neutrinos

$$\Gamma(H \to N_R N_R) = \frac{m_{N_R}^2 m_H}{16 \pi v_{\text{B-L}}^2} \left(1 - \frac{4m_{N_R}^2}{m_H^2}\right)^{3/2} (\cos\alpha)^2$$
(A.4)



Figure A.1: Plot of heavy scalar boson decay width as a function of scalar mixing angle $\cos \alpha$ for different values of m_H [168].

Decay of heavy scalar into the SM like Higgs

$$\Gamma(H \to hh) = \frac{\lambda_{Hhh}^2}{32 \ \pi \ m_H} \sqrt{1 - \frac{4m_h^2}{m_H^2}}$$
(A.5)

Figure. A.1 shows the dependence of total decay width of the heavy scalar boson $\Gamma_{\rm H}^{tot}$ on the scalar mixing $\cos \alpha$ for different values of m_H . For higher m_H , the decay-width becomes larger for large mixing. This plot also shows that for the limiting case when $\cos \alpha \rightarrow 1.0$, i.e, without mixing between the scalar bosons, $\Gamma_{\rm H}^{tot} \rightarrow 0$ and hence it is completely de-coupled from the SM.

Appendix B

Detailed calculation of w(s)

Let ϕ be the scattering angle between incoming DM particles then w(s) can be defined [161, 170] as

$$w(s) = \frac{1}{32\pi} \sqrt{\frac{s - 4m_{final}^2}{s}} \int \frac{d\cos\phi}{2} \sum_{\text{all possible channels}} |\mathcal{M}|^2.$$
(B.1)

The function $|\mathcal{M}|^2$ contains not only interaction part, but also contains the kinematical part. Considering the processes as in eq. (4.18) we can write

$$w(s)_{b,\tau,W,Z} = \left[\frac{\sin^{2} \alpha \cos^{2} \alpha}{4} \left(4y_{n_{3}}^{2}(s-4m_{N_{R}}^{2})\right)\right] \times \left[\frac{1}{(s-m_{h}^{2})^{2}+\Gamma_{h}^{2}m_{h}^{2}} + \frac{1}{(s-m_{H}^{2})^{2}+\Gamma_{H}^{2}m_{H}^{2}} -2\frac{(s-m_{h}^{2})(s-m_{H}^{2})+m_{h}m_{H}\Gamma_{h}\Gamma_{H}}{((s-m_{h}^{2})^{2}+\Gamma_{h}^{2}m_{h}^{2})((s-m_{H}^{2})^{2}+\Gamma_{H}^{2}m_{H}^{2})}\right] \times \left[\left\{\frac{1}{8\pi}\sqrt{\frac{s-m_{b}^{2}}{s}} 4y_{b}^{2}\left(\frac{s}{4}-m_{b}^{2}\right)3\right\} + \left\{\frac{1}{8\pi}\sqrt{\frac{s-m_{\tau}^{2}}{s}} 4y_{\tau}^{2}\left(\frac{s}{4}-m_{\tau}^{2}\right)\right\} + \left\{\frac{1}{8\pi}\sqrt{\frac{s-m_{W}^{2}}{s}}\left(\frac{2m_{W}^{2}}{v}\left(s+\frac{1}{2m_{W}^{4}}\left(\frac{s}{2}-m_{W}^{2}\right)\right)\right)\right\}\right].$$

$$\left. + \left\{\frac{1}{8\pi}\sqrt{\frac{s-m_{Z}^{2}}{s}}\left(\frac{m_{Z}^{2}}{v}\left(s+\frac{1}{2m_{Z}^{4}}\left(\frac{s}{2}-m_{Z}^{2}\right)\right)\right)\right\}\right].$$
(B.2)

In this expression second line is the propagator function which includes both hand H. Third line shows decay cross section to $b\bar{b}$ and $\tau^+\tau^-$, whereas, fourth and fifth line is decay cross section to W^+W^- and ZZ respectively. In addition, we have also considered the annihilation into the SM-like Higgs bosons, for which $w(s)_h$ is given by,

$$w(s)_{h} = \left\{ \frac{1}{16\pi} \left[4y_{n_{3}}^{2} (s - 4m_{N_{R}}^{2}) \right] \sqrt{\frac{s - m_{h}^{2}}{s}} \\ \left(\left(\frac{\sin\alpha}{\sqrt{2}} \right)^{2} \frac{\lambda_{hhh}^{2}}{(s - m_{h}^{2})^{2} + \Gamma_{h}^{2} m_{h}^{2}} + \left(\frac{\cos\alpha}{\sqrt{2}} \right)^{2} \frac{\lambda_{Hhh}^{2}}{(s - m_{H}^{2})^{2} + \Gamma_{H}^{2} m_{H}^{2}} - \frac{\sin\alpha \cos\alpha \lambda_{hhh} \lambda_{Hhh} \left\{ (s - m_{h}^{2})(s - m_{H}^{2}) + m_{h} m_{H} \Gamma_{h} \Gamma_{H} \right\}}{((s - m_{h}^{2})^{2} + \Gamma_{h}^{2} m_{h}^{2}) ((s - m_{H}^{2})^{2} + \Gamma_{H}^{2} m_{H}^{2})} \right) \right\}$$
(B.3)

where, λ_{hhh} and λ_{Hhh} are calculated by expanding the Higgs potential part,

$$\lambda_{hhH} = 3\lambda_1 v \left(\cos^2 \alpha \sin \alpha\right) + 3\lambda_2 v_{\text{B-L}} \left(\cos \alpha \sin^2 \alpha\right) \\ + \frac{1}{8}\lambda_3 \left\{ v_{\text{B-L}} \left(\cos \alpha + 3\cos(3\alpha)\right) + v \left(\sin \alpha - 3\sin(3\alpha)\right) \right\}, \\ \lambda_{hhh} = \frac{\lambda_1}{4} v \left(3\cos \alpha + \cos(3\alpha)\right) + \frac{\lambda_2}{4} v_{\text{B-L}} \left(-3\sin \alpha + \sin(3\alpha)\right) \\ + \frac{\lambda_3}{8} \left\{ v \left(\cos \alpha - \cos(3\alpha)\right) - v_{\text{B-L}} \left(\sin \alpha + \sin(3\alpha)\right) \right\}.$$
(B.4)

Finally, $w(s) = w(s)_{b,\tau,W,Z} + w(s)_h$.

Appendix C

Loop functions involved in $\langle \sigma v \rangle_{\gamma\gamma}$

The loop functions involved in Higgs to di-photon process [171, 172] are depicted as:

$$F_t(\tau) = -2\tau [1 + (1 - \tau)f(\tau)] ,$$

$$F_W(\tau) = 2 + 3\tau + 3\tau (2 - \tau)f(\tau) ,$$

where, $\tau_i = 4m_i^2/m_{h,H}^2$ (i = W, t) and

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

For, $m_h = 125$ GeV the loop-functions becomes,

$$F_t(\tau_t) = 1.83$$
, $F_W(\tau_W) = -8.32$.

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Publications attached with the thesis

- Triplet-singlet extension of the MSSM with a 125 GeV Higgs boson and dark matter,
 Tanushree Basak and Subhendra Mohanty, Phys. Rev. D, 86, 075031 (2012).
- Constraining Minimal U(1)_{B-L} model from Dark Matter Observations, Tanushree Basak and Tanmoy Mondal, Phys. Rev. D, 89, 063527 (2014).