The recent observed anomalies \cite{1,2,3} in the cosmic rays have given a lot of excitement. It has been shown that there is a clean excess of absolute positron flux in the cosmic rays at an energy \( E \gtrsim 50 \text{ GeV} \) \cite{4}, even if the propagation uncertainty \cite{5} in the secondary positron flux is added to the Galactic background. This leaves enough motivation for considering particle physics motivated dark matter (DM) models, such as annihilation \cite{6,7} or decay \cite{7,8} of DM, as the origin of positron excess in the cosmic rays\footnote{For astrophysical origins, see Ref. \cite{3} and references therein.}. However the origin of the DM, being interpreted as a long lived particle, goes beyond the standard model (SM). Moreover, the origin of observed matter antimatter asymmetry and the origin of inflation are also two crucial phenomena which require the physics beyond the SM.

In particular once cosmic inflation occurs, then after inflation it must pave the way to excite not only the observed SM quanta but also the DM. This can be achieved minimally if the inflaton, \( \phi \), itself carries the SM charges as in the case of the following examples which relies on the minimal supersymmetric SM (MSSM) setup \cite{10,11,12}. In the MSSM case the inflaton decays into the SM quarks and leptons and through thermal scatterings the lightest supersymmetric DM particles were created which are absolutely stable. On the other hand if inflation does not happen in the observable sector, for instance if it belongs to a hidden sector, as in the case of plethora of examples \cite{13}, then the onus will be to explain how to generate the desired degrees of freedom, i.e. SM baryons and dark matter abundance.

The aim of this paper is to illustrate an example where lepton asymmetry \cite{14} and DM abundance are generated right above the electroweak scale, i.e., \( \mathcal{O}(100) \text{ GeV} \). Our building block of beyond the SM physics is based on a crucial observation of cosmic ray anomalies \cite{2}, which may unravel the mystery of these issues in a unifying framework where the DM particle itself carries a net \( B-L \) asymmetry and decays very slowly to the SM particles. Note that previous attempts \cite{15,16,17} were made to unify dark matter and baryogenesis, but it is more challenging to address why the DM annihilates/decays primarily into leptons and anti-leptons in a unifying framework, thus explaining the observed cosmic ray anomalies at PAMELA \cite{1} and Fermi \cite{3}.

For the purpose of illustration, let us augment the SM by adding a new \( U(1)_{B-L} \) gauge symmetry, and without supersymmetry. The anomaly free gauged \( B-L \) symmetry then naturally accommodates a new fermionic dark matter \( N_L(1,0,-1) \), where the quantum numbers inside the parenthesis shows the transformation properties of \( N_L \) under the gauge group \( SU(2)_L \times U(1)_Y \times U(1)_{B-L} \). At a high scale the \( U(1)_{B-L} \) gauge symmetry is broken by a scalar field and give mass \( M_N = F v_{B-L} \) to \( N_L = (1/\sqrt{2})(N_L + N_R) \), where \( v_{B-L} \) is the vacuum expectation value (vev) of the \( U(1)_{B-L} \) breaking scalar field which carries \( B-L \) charges by two units and \( F \) is the coupling between \( B-L \) scalar field and \( N_L \).

In the broken phase of the \( U(1)_{B-L} \) gauge symmetry the Lagrangian involving the interactions of \( N_L \), and the new massive charged scalars, \( \eta^- (1,-2,0) \) and \( \chi^- (1,-2,-2) \), can be separately written in terms of \( B-L \) conserving and \( B-L \) violating parts. The \( B-L \) violating of the Lagrangian is given by:

\[
\mathcal{L}_{\Delta(B-L) \neq 0} = \frac{1}{2} (M_N)_{\alpha \beta} (\bar{N}_{\alpha L} \eta^{\dagger} N_{\beta L} + m^2 \bar{N}_{\alpha L} \chi^{\dagger} + h.c.), \quad (1)
\]

with \( m^2 = \mu' v_{B-L} \), while the relevant \( B-L \) conserving part of the Lagrangian is given by:

\[
\mathcal{L}_{\Delta(B-L)=0} \supset M_{\alpha}^{2} \eta^{\dagger} \eta + M_{\chi}^{2} \chi^{\dagger} \chi + \mu \bar{H}_1 H_2
\]

\[
+ h_{\alpha \beta} \eta^{\dagger} (\bar{N}_{\alpha L} \ell_{\beta R} + f_{\alpha \beta} \chi^{\dagger} \ell_{\alpha L} \ell_{\beta L} + h.c.)(2)
\]

where the indices \( \alpha, \beta = e, \mu, \tau \) represent the flavor basis of the SM fermion fields. In equation (2), \( \ell_{L}(2,-1,-1) \) and \( \ell_{R}(1,-2,-1) \) represent the lepton doublet and singlet respectively, while \( H_1(2,1,0) \), \( H_2(2,1,0) \) are two doublet Higgses which couple to up and down sector of SM fermions.

An important point to note is that the mass term, \( m^2 \eta^{\dagger} \eta \), which violates \( B-L \) by two units, gives rise a mixing between \( \eta \) and \( \chi \). Due to this mixing the Majorana fermion \( N \) will be an unstable leptonic DM. Therefore, the decay products of \( N \) are only SM leptons/antileptons. If the lifetime of \( N \) is about \( \mathcal{O}(10^{25}) \) s or so, then the decay products of \( N \) can naturally account for the observed \( e^\pm \) excesses at PAMELA and Fermi. This observation will place non-trivial constraint, such as \( m < M_{\eta}, M_{\chi} \). Further we note that the only part of the Lagrangian is given by:

\[
\mathcal{L}_{\Delta(B-L) \neq 0} = \frac{1}{2} (M_N)_{\alpha \beta} (\bar{N}_{\alpha L} \eta^{\dagger} N_{\beta L} + m^2 \bar{N}_{\alpha L} \chi^{\dagger} + h.c.), \quad (1)
\]

with \( m^2 = \mu' v_{B-L} \), while the relevant \( B-L \) conserving part of the Lagrangian is given by:

\[
\mathcal{L}_{\Delta(B-L)=0} \supset M_{\alpha}^{2} \eta^{\dagger} \eta + M_{\chi}^{2} \chi^{\dagger} \chi + \mu \bar{H}_1 H_2
\]

\[
+ h_{\alpha \beta} \eta^{\dagger} (\bar{N}_{\alpha L} \ell_{\beta R} + f_{\alpha \beta} \chi^{\dagger} \ell_{\alpha L} \ell_{\beta L} + h.c.)(2)
\]

where the indices \( \alpha, \beta = e, \mu, \tau \) represent the flavor basis of the SM fermion fields. In equation (2), \( \ell_{L}(2,-1,-1) \) and \( \ell_{R}(1,-2,-1) \) represent the lepton doublet and singlet respectively, while \( H_1(2,1,0) \), \( H_2(2,1,0) \) are two doublet Higgses which couple to up and down sector of SM fermions.

An important point to note is that the mass term, \( m^2 \eta^{\dagger} \eta \), which violates \( B-L \) by two units, gives rise a mixing between \( \eta \) and \( \chi \). Due to this mixing the Majorana fermion \( N \) will be an unstable leptonic DM. Therefore, the decay products of \( N \) are only SM leptons/antileptons. If the lifetime of \( N \) is about \( \mathcal{O}(10^{25}) \) s or so, then the decay products of \( N \) can naturally account for the observed \( e^\pm \) excesses at PAMELA and Fermi. This observation will place non-trivial constraint, such as \( m < M_{\eta}, M_{\chi} \). Further we note that the only
coupling “\(h\)” is responsible for the production of a net lepton asymmetry and DM in the early universe, while the three body decay of DM through the same coupling at current epoch gives rise to the observed anomalies at PAMELA and Fermi.

**Baryon asymmetry:**

Let us now consider how can we explain the observed matter-anti-matter asymmetry in our setup. A natural possibility is to generate a lepton asymmetry from the out-of-equilibrium decay of \(\eta^-\) field. Similar to Dirac leptogenesis [13], here we will argue that the required baryon asymmetry can be generated from a conserved \(B - L\) number. In order to generate the baryon asymmetry we need following three steps:

1. At first the CP-violating out-of-equilibrium decay of \(\eta^-\) must generate an equal and opposite \(B - L\) asymmetry between \(N_L\) and \(\ell_R\). These two asymmetries should not equilibrate above the electron weak (EW) phase transition.

2. Above the EW-phase transition the \(B - L\) asymmetry stored in \(\ell_R\) gets transferred to \(\ell_L\), while keeping an equal and opposite \(B - L\) asymmetry in \(N_L\).

3. The \(B - L\) asymmetry stored in \(\ell_L\) then gets converted to a net baryon asymmetry in the presence of \(SU(2)_L\) sphalerons, while keeping the \(B - L\) asymmetry stored in \(N_L\) intact.

Note that all three steps happen right above the EW scale. Since \(\eta^-\) is neutral under \(B - L\), it can decay to a pair of lepton (\(\ell_R\)) and antilepton (\(\overline{N}_L\)). Note that the decay of \(\eta^-\) cannot produce any lepton asymmetry since its decay does not violate any lepton number. However, if there are at least two \(\eta^-\) fields, say \(\eta_1^\pm\) and \(\eta_2^\pm\), then there can be CP violation in the decay of \(\eta^-\) fields. In their mass basis, spanned by \(\psi_1^\pm\) and \(\psi_2^\pm\), the lightest \(\psi^\pm\) field, say \(\psi_1^\pm\), can generate a net CP asymmetry through the interference of tree level and self energy correction diagram [20]. The CP asymmetry is then given by

\[
\epsilon_1 = \frac{\text{Im} \left[ (\mu_1 \mu_2^*)^2 \sum_{ij} h_{ij}^L h_{ij}^{L*} \right]}{16\pi^2 (M_{\eta_2}^2 - M_{\eta_1}^2)} \left( \frac{M_{\eta_1}}{\Gamma_{\eta_1}} \right),
\]

where

\[
\Gamma_{\eta_1} = \frac{1}{8\pi M_{\eta_2}} \left( \mu_1 \mu_2^2 + M_{\eta_1} M_{\eta_2} \sum_{i,j} h_{ij}^L h_{ij}^{L*} \right).
\]

Now assuming

\[
\frac{M_{\eta_1}^2 M_{\eta_2}^2}{(M_{\eta_2}^2 - M_{\eta_1}^2)} = O(1), \quad \frac{\mu_1 \mu_2}{M_{\eta_1} M_{\eta_2}} = O(1)
\]

and

\[
h_{ij}^L \approx h_{ij}^\eta = O(10^{-2}),
\]

we get from Eqns. (3), (4) and (5) the CP asymmetry \(\epsilon_1 \approx 10^{-5}\). Due to the CP violation the decay of \(\psi_1^\pm\) generates an equal and opposite \(B - L\) asymmetry between \(N_L\) and \(\ell_R\). Since the interaction between \(N_L\) and \(\ell_R\) through the coupling ‘\(h\)’ is already gone out-of-equilibrium, the asymmetries between them don’t equilibrate any more at the required scale of \(O(100)\) GeV. On the other hand, the lepton number conserving process: \(\ell_R \ell_R \leftrightarrow \ell_L \ell_L\), mediated via the SM Higgs, remains in thermal equilibrium above the electroweak phase transition. As a result the \(B - L\) asymmetry stored in \(\ell_R\) gets transferred to \(\ell_L\) through this L-number conserving process, while leaving an equal and opposite \(B - L\) asymmetry in \(N_L\). The transportation of \(B - L\) asymmetry from \(\ell_R\) to \(\ell_L\) can be understood as follows. Let us define the chemical potential associated with the \(\ell_R\) field as \(\mu_{\ell_R} = \mu_0 + \mu_{BL}\), where \(\mu_{BL}\) is the chemical potential contributing to \(B - L\) asymmetry and \(\mu_0\) is independent of \(B - L\). Hence at equilibrium we have the chemical potential associated with \(\ell_L\) is given by \(\mu_{\ell_L} = \mu_{\ell_R} + \mu_H = \mu_{BL} + \mu_0 + \mu_H\). Thus we see that the same chemical potential is associated with \(\ell_L\) as of \(\ell_R\). Therefore, the net \(B - L\) asymmetry stored in \(\ell_L\) can be passed on to \(\ell_L\). Since the \(SU(2)_L\) sphalerons are in thermal equilibrium at a scale above 100 GeV, the \(B - L\) asymmetry stored in \(\ell_L\) can be converted to a net baryon asymmetry, while an equal and opposite \(B - L\) asymmetry will remain in \(N_L\). The two asymmetries will equilibrate when \(N_L\) will decay through the \(B - L\) violating process. The net \(B - L\) asymmetry thus produced can be given as:

\[
\eta_{B-L} = \frac{3}{4} B_\eta \epsilon_1 \frac{T_R}{m_\phi},
\]

where \(B_\eta\) is the inflaton branching ratio, which is of order, \(O(1)\), and \(T_R \gtrsim 100 \text{ GeV}\) is the reheating temperature of the universe and \(m_\phi\) is the inflaton mass.

The conversion of lepton asymmetry to the baryon asymmetry is obtained by \(\eta_B = (28/79)\eta_{B-L}\). For \(T_R/m_\phi \approx 10^{-2}\) and \(\epsilon_1 \approx 10^{-5}\), we can achieve the observed baryon asymmetry \(\eta_B \approx O(10^{-10})\). A crucial point to note here is that the lepton asymmetry is virtually created by a non-thermal decay of the inflaton decay products, \(\eta\) and \(\chi\), which we will discuss below.

**Dark matter abundance:**

Next we discuss the number density of the lightest \(N_L\) to check if it satisfies the observed DM abundance. It turns out that a large abundance of \(N_L\) will be produced non-thermally by the decay of \(\eta\), which is also non-thermally produced by the inflaton decay.
As we will argue below, the annihilation cross section of $N_L$ is larger than the canonical one, $\langle |v| \rangle \sim 3 \times 10^{-26}$ cm$^3$s$^{-1}$. In this case the final abundance of the thermal component is much smaller than the observational value,

$$Y_{\text{DM}} \equiv \frac{n_{\text{DM}}}{s} = 4 \times 10^{-13} \left( \frac{1 \text{ TeV}}{M_{\text{DM}}} \right) \left( \frac{\Omega_{\text{DM}} h^2}{0.11} \right), \quad (7)$$

where $M_{\text{DM}}$ is the DM mass, $\Omega_{\text{DM}}$ is the density parameter of the DM with $h$ being the normalized Hubble constant.

Since we consider a case when the inflaton, $\phi$, decays well after the standard freeze-out epoch of the thermally-produced $N_L$, i.e., at cosmic temperature $T < M_N/25$. Then the yield value of $N_L$ is estimated by

$$Y_{N_L} = \frac{n_{N_L}}{s} \approx 3 \times 10^{-3} B_\eta \frac{T_R}{m_\phi}, \quad (8)$$

For $T_R/m_\phi \sim 10^{-4}$, as required by the lepton asymmetry, we find that the relic abundance of $N_L$ is given by $Y_{N_L} \approx 10^{-5}$, which is much larger than the observed DM abundance $7$.

However thanks to the gauge coupling of $N_L$, such that $N_L$ can now annihilate into the SM fermions through the exchange of $Z_{B-L}$ gauge boson after its non-thermal production. We can then obtain the final abundance of $N_L$ by solving the Boltzmann equations:

$$\frac{dn_L}{dt} + 3n_L H = -\langle |v| \rangle n^2_{N_L} + \Gamma n_L,$$

where $\langle |v| \rangle \approx (1/4\pi) M_N^2/v_{B-L}^4$, and we have omitted the production term from the thermal bath, $+(\langle |v| \rangle n^2_{N_L})_{\text{eq}}$ in the right-hand side of the second line. Then we approximately obtain,

$$Y_{N_L} \approx \frac{3H}{\langle |v| \rangle} s. \quad (10)$$

The right-hand side of Eq. (10) is $\propto 1/T_R$, which means that the late-time decay induces a larger freeze-out value. In Fig. 1 we illustrate such a non-thermal production and/or further annihilation mechanism.

By equating $Y_{N_L} \approx Y_{\text{DM}}$ we get a constraint on the $B-L$ breaking scale with satisfying the observational DM density to be

$$v_{B-L} \approx 5 \times 10^3 \text{ GeV} \left( \frac{\Omega_{\text{DM}} h^2}{0.11} \right)^{1/4} \left( \frac{M_N}{3 \text{ TeV}} \right)^{1/4} \times \left( \frac{T_R}{100 \text{ GeV}} \right)^{1/4}. \quad (11)$$

This gives $\langle |v| \rangle \sim O(10^{-25}) - O(10^{-24})$ cm$^3$s$^{-1}$ for $M_N \sim O(\text{TeV})$, which is larger than the canonical annihilation cross section and makes the thermal component sub dominant. In turn, we understand this mechanism intuitively by using Fig. 1 as follows. If we specify a $U(1)_{B-L}$ breaking scale ($\sim \text{TeV}$ in this model) or an annihilation cross section, the model necessarily has a crossing point between the diagonal and the horizontal lines at $x = M_{\text{DM}}/T_R$ shown in Fig. 1 which gives the right observational abundance of DM and its production epoch. In the current model, we demand the cross section to be larger than the canonical value and $T_R$ to be larger than 100 GeV.

Note that the annihilation process $N_L N_L \rightarrow \bar{f} f$ is a $B-L$ number conserving process and therefore does not transfer any $B-L$ asymmetry to the SM fermions. As a result the $B-L$ asymmetry produced via the decay of $\eta^-$ will survive until far below the electro-weak phase transition.

Since $B-L$ is already broken, the lightest $N_L$ is no more stable. It will decay through the three body process: $N_L \rightarrow e^{-} \mu^{+} \bar{\nu}_{\mu} \bar{\nu}_{\tau}$, with $\beta \neq \gamma$, through the mixing of $\eta$ and $\chi$. Since the coupling of $\chi$ to two lepton doublets is antisymmetric, i.e., $\beta \neq \gamma$, the decay of $N_L$ is not necessarily to be flavour conserving. In particular the decay mode: $N_L \rightarrow \tau_{R} \tau_{L} \bar{\nu}_{e} \bar{\nu}_{\mu}$, violates $L_{e}$ ($L_{\mu}$) by one unit while it violates $L = L_{e} + L_{\mu} + L_{\tau}$ by two units.
In the mass basis of $N_L$ the life time can be estimated to be

$$\tau_N = 2.0 \times 10^{25}s \left( \frac{10^{-2}}{\hbar} \right)^2 \left( \frac{10^{-7}}{f} \right)^2 \left( \frac{50 \text{ GeV}}{m} \right)^4 \left( \frac{m_\phi}{10^6 \text{ GeV}} \right)^8 \left( \frac{1 \text{ TeV}}{M_{N_L}} \right)^5,$$  \hspace{1cm} (12)

where we assume that $M_\eta \simeq M_\kappa \approx m_\phi$ in order to get a lower limit on the lifetime of $N_L$. The prolonged life time of $N_L$ may explain the current cosmic ray anomalies observed by PAMELA \cite{1} and Fermi \cite{3}. The electron and positron energy spectrum can be estimated by using the same set-up as in Ref. \cite{8}. In Figs. 2 and 3 we have shown the integrated $e^\pm$ fluxes in a typical decay mode: $N_L \rightarrow \tau^-\tau^+\bar{\nu}$ up to the maximum available energy $M_N/2$ for $\tau_N = 4.0 \times 10^{25}$ secs. From there it can be seen that the decay of $N_L$ can nicely explain the observed positron excess at PAMELA and $e^\pm$ excesses at Fermi. While doing so we assume that the branching fraction in the decay of $N - L$ to $\tau^-\tau^+\bar{\nu}$ is significantly larger than the other viable decay modes: $N_L \rightarrow \mu^-\mu^+\bar{\nu}$ and $N_L \rightarrow e^-e^+\bar{\nu}$. However, if the decay rate of $N_L \rightarrow \mu^-\mu^+\bar{\nu}$ is comparable to $N_L \rightarrow \tau^-\tau^+\bar{\nu}$ then it can explain the observed anomalies at PAMELA and Fermi, while the decay mode: $N_L \rightarrow e^-e^+\bar{\nu}$ produces larger $e^+e^-$ fluxes at Fermi and therefore unfavorable.

Another potential signature of this scenario is the emission of energetic neutrinos from the Galactic center \cite{23} which can be checked by future experiments such as IceCube DeepCore \cite{26} and KM3NeT \cite{27}. We will come back to this issue in a future work \cite{28}.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig2.pdf}
\caption{Positron excess from $N_L \rightarrow \tau^-\tau^+\bar{\nu}$ with $M_N = 3$ TeV. The fragmentation function has been calculated using PYTHIA \cite{24}.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig3.pdf}
\caption{$e^\pm$ excess from $N_L \rightarrow \tau^-\tau^+\bar{\nu}$ with $M_N = 3$ TeV. The fragmentation function has been calculated using PYTHIA \cite{24}.}
\end{figure}

Inflation and reheating:
So far we have not discussed anything about inflation sector. Let us now consider a hidden sector inflaton $\phi(1, 0, 0)$ which excites the observed DM abundance and SM leptons during the process of reheating which occurs after the end of inflation. Let us consider a simple toy model of inflation, where the inflaton potential admits a point of inflation, see \cite{10, 11, 15, 21}:

$$V(\phi) \sim \frac{m_\phi^2}{2} \phi^2 - \frac{A\kappa}{6\sqrt{3}} \phi^3 + \frac{\kappa^2}{12} \phi^4,$$  \hspace{1cm} (13)

where $A \approx 4m_\phi$ and $\kappa \sim 10^{-10}$. Inflation can happen near a point of inflation given by $\phi_0 \sim \sqrt{3}m_\phi/\kappa \sim 10^{16}$ GeV with an Hubble expansion rate, $H_{inf} \sim (m_\phi^2/\kappa M_P) \sim 10^4$ GeV. The amplitude of the density perturbations will be given by: $\delta_H \approx (1/5\pi)(H_{inf}^2/\phi) \sim (\kappa^2 M_P/3m_\phi) N^2 \sim 10^{-5}$, where the number of e-folds is given by: $N^2 \sim 10^3$ \cite{22}. One of the dynamical properties of an inflation point inflation is that the spectral tilt can be matched in a desired observable range: $0.92 < n_s < 1.0$ for the above parameters \cite{10, 11, 15, 21}.

The inflaton decays into heavy charged scalars $\eta$ and $\chi$. Let $B_\eta$ be the branching fraction in the decay of $\phi$ to $\eta^\pm$ and $B_\chi$ be the branching fraction in the decay of $\phi$ to $\chi^\pm$. As we discussed above the charged scalars $\eta^\pm$ and $\chi^\pm$ couple to the SM degrees of freedom. The reheating occurs when the inflaton begins oscillations. The largest decay rate happens for the largest amplitude of oscillations, for instance when $\langle \phi \rangle \sim \phi_0$, see for instance \cite{24}. As a result the Universe gets reheated up to a desired temperature:

$$T_R \sim 0.1 \sqrt{\Gamma_\phi M_P} = 1.2 \times 10^2 \text{GeV} \left( \frac{g}{10^{-17}} \right) \left( \frac{m_\phi}{10^6 \text{ GeV}} \right)^{1/2},$$  \hspace{1cm} (14)
above the electro-weak scale to facilitate a successful baryogenesis, where \( \Gamma_\phi = \left( g^2/8\pi \right) \left( (\phi/m_\phi)^2 \right) m_\phi \) is the decay rate of \( \phi \) with \( g \) is the quartic coupling; \( g^2(\eta^+ \eta^- + \chi^\dagger \chi) \). A small decay rate of inflaton to \( \eta^\pm \) and \( \chi^\pm \) ensures the optimal temperature just right above the electro-weak phase transition.

To summarize, we have explored a simple model of lepto-philic universe where DM carries a net \( B-L \) asymmetry and decays only into the SM leptons can explain the observed positron excess at PAMELA and \( e^\pm \) excesses at Fermi. These anomalous observations at PAMELA and Fermi may indirectly probe the common origin of mysterious DM and baryon asymmetry as we have shown. The baryon asymmetry in our model is created via lepton conserving leptogenesis mechanism which gets converted into baryon asymmetry via the electro-weak sphalerons. Before closing we note that the model explained here can be embedded within supersymmetry without further challenges, neither the leptogenesis nor the DM mechanisms will alter, the parameters and the observed values of the inflationary perturbations and the tilt in the spectrum would remain so. The details will be published elsewhere [28].

Acknowledgement: The authors are supported by the European Union through the Marie Curie Research and Training Network “UniverseNet” (MRTN-CT-2006-035863) and STFC grant, PP/D000394/1. NS would like to thank Utpal Sarkar for useful discussions.

[1] O. Adriani et al., [arXiv:0810.4995 [astro-ph]]
[2] J. Chang et al., Nature 456, 362 (2008); S. Tori et al., [arXiv:0909.0760 [astro-ph]]; F. Aharonian et al., [HESS Collaboration], [arXiv:0905.0105 [astro-ph.HE]]; J. J. Beatty et al., Phys. Rev. Lett. 93 (2004) 241102; M. Aguilar et al. [AMS-01 Collaboration], Phys. Lett. B 646, 145 (2007).
[3] A. A. Abdo et al., [Fermi LAT Collaboration], arXiv:0905.0025 [astro-ph.HE].
[4] C. Balazs, N. Sahu and A. Mazumdar, JCAP 0907, 039 (2009) arXiv:0905.4302 [hep-ph].
[5] I. V. Moskalenko and A. W. Strong, Astrophys. J. 493, 694 (1998); T. Delahaye, R. Lineros, F. Donato, N. Fornengo and P. Salati, Phys. Rev. D 77, 063527 (2008); E. A. Baltz and J. Edsjo, Phys. Rev. D 59, 023511 (1999).
[6] L. Bergstrom, J. Edsjo and G. Zaharijas, arXiv:0905.0333 [astro-ph.HE]; D. Hooper and P. Salati, Phys. Rev. D 77, 060990 (2008); T. Delahaye, R. Lineros, F. Donato, N. Fornengo and P. Salati, Phys. Rev. D 77, 063527 (2008); E. A. Baltz and J. Edsjo, Phys. Rev. D 59, 023511 (1999).
[7] K. Hamaguchi, K. Nakaji and E. Nakamura, arXiv:0905.1574 [hep-ph].
[8] K. Ishiwata, S. Matsumoto and T. Moroi, JHEP 0905, 110 (2009) arXiv:0903.0242 [hep-ph]; A. Ibarra and D. Tran, JCAP 0902, 021 (2009); S. Shirai, F. Takahashi and T. T. Yanagida, arXiv:0905.3255 [hep-ph]; A. Arvanitaki, S. Dimopoulos, S. Dubovsky, P. W. Graham, R. Harnik and S. Rajendran, arXiv:0904.2789 [hep-ph]; C. H. Chen, C. Q. Geng and D. V. Zhuridov, arXiv:0905.0652 [hep-ph]; N. Okada and T. Yamada, arXiv:0906.2801 [hep-ph].
[9] D. Hooper, P. Blasi, and P. Dario Serpico, Journal of Cosmology and Astro-Particle Physics 1, 25 (2009); H. Yukawa, M. D. Kistler, and T. Stanek, arXiv:0810.2784 [astro-ph,G.A]; S. Profumo, arXiv:0812.4577; K. Ioka, arXiv:0812.4851; E. Borriello, A. Cuoco and G. Miele, arXiv:0903.1852 [astro-ph.GA]; P. Blasi, arXiv:0903.2704; P. Blasi and P. D. Serpico, arXiv:0904.0871; N. Kawamura, K. Ioka and M. M. Nojiri, arXiv:0903.3782 [astro-ph.HE]; Y. Fujita, K. Kohri, R. Yamazaki and K. Ioka, arXiv:0906.5298 [astro-ph.HE].
[10] R. Allahverdi, K. Enqvist, J. Garcia-Bellido and A. Mazumdar, Phys. Rev. Lett. 97, 191304 (2006) R. Allahverdi, K. Enqvist, J. Garcia-Bellido, A. Jokinen and A. Mazumdar, JCAP 0706, 019 (2007)
[11] R. Allahverdi, A. Kusenko and A. Mazumdar, JCAP 0707, 018 (2007)
[12] R. Allahverdi, B. Dutta and A. Mazumdar, Phys. Rev. D 75, 075018 (2007) R. Allahverdi, B. Dutta and A. Mazumdar, Phys. Rev. Lett. 99, 261301 (2007)
[13] A.D. Linde, Particle Physics and Inflationary Cosmology (Harwood Academic Publishers, Chur, Switzerland 1990).
[14] M. Fukugita and T. Yanagida, Phys. Lett. B 174 45.1986.
[15] K. Kohri, A. Mazumdar and N. Sahu, arXiv:0905.1625 [hep-ph].
[16] M. Aoki, S. Kanemura and O. Seto, Phys. Rev. Lett. 102, 051805 (2009); K. S. Babu and E. Ma, Int. J. Mod. Phys. A 23, 1813 (2008); T. Hambye, K. Kannike, E. Ma and M. Raidal, Phys. Rev. D 75, 095003 (2007); E. Ma, Mod. Phys. Lett. A 21, 1777 (2006); P. H. Gu and U. Sarkar, Phys. Rev. D 77, 105031 (2008); N. Sahu and U. A. Yajnik, Phys. Lett. B 635, 11 (2006) arXiv:hep-ph/0509285.
[17] N. Sahu and U. Sarkar, Phys. Rev. D 76, 045014 (2007); J. McDonald, N. Sahu and U. Sarkar, JCAP 0804, 037 (2008); N. Sahu and U. Sarkar, Phys. Rev. D 78, 115013 (2008).
[18] M. Frigerio, T. Hambye and E. Ma, JCAP 0609, 009 (2006) arXiv:hep-ph/0603123.
[19] In case of “Dirac Leptogenesis” there is no need of Lepton number violation. See for instance K. Dick, M. Lindner, M. Ratz and D. Wright, Phys. Rev. Lett. 84, 4039 (2000).
[20] E. Ma and U. Sarkar, Phys. Rev. Lett. 80, 5716 (1998) arXiv:hep-ph/9802445.
[21] J. C. Bueno Sanchez, K. Dimopoulos and D. H. Lyth, JCAP 0701 (2007) 015; R. Allahverdi, B. Dutta and A. Mazumdar, Phys. Rev. D 78, 063507 (2008); R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0610069.
[22] C. P. Burgess, R. Easther, A. Mazumdar, D. F. Mota and T. Multamaki, JHEP 0505, 067 (2005).
[23] R. Allahverdi and A. Mazumdar, JCAP 0610, 008 (2006).
[24] T. Sjostrand, S. Mrenna and P. Skands, JHEP 0605, 026.
[25] J. Hisano, M. Kawasaki, K. Kohri and K. Nakayama, Phys. Rev. D79 (2009) 043516; J. Hisano, K. Nakayama and M. J. S. Yang, arXiv:0905.2075 [hep-ph]; J. Liu, P. f. Yin and S. h. Zhu, arXiv:0812.0964 [astro-ph].
[26] D. F. Cowen [IceCube Collaboration], J. Phys. Conf. Ser. 110, 062005 (2008).
[27] A. Kappes and f. t. K. Consortium, arXiv:0711.0563 [astro-ph].
[28] K. Kohri, A. Mazumdar, N. Sahu and P. Stephens, under preparation.