Supersymmetry (SUSY) elegantly solves the quadratic ultra-violet (UV) sensitivity of the Higgs mass via the introduction of particles (superpartners) with opposite statistics to each Standard Model (SM) particle. Stability of the proton naturally leads to the introduction of R-parity under which all superpartners are odd while the lightest supersymmetric content is even. Thus sparticles can only be created in pairs at colliders and the lightest supersymmetric particle (LSP) is stable, providing an interesting dark matter (DM) candidate.

In the minimal supersymmetric extension of the SM a natural and well-studied DM candidate is the lightest neutralino $\tilde{\chi}_1^0$, a linear combination of the wino $\tilde{W}$, bino $\tilde{B}$ and Higgsinos $\tilde{h}_u$ and $\tilde{h}_d$ superpartners. Given its weak couplings and for masses of order the EW scale, the lightest neutralino can provide the well-known "WIMP miracle" in which the total amount of DM relic density is naturally obtained. Despite its appealing properties, neutralino DM in the MSSM is being pushed toward corners of parameter space, in particular due to the lack of positive signals in direct detection experiments that probe spin-independent (SI) [1], [2] and spin-dependent (SD) [3] scattering of DM particles off of nuclei target. Furthermore, the absence of discovery of superpartners at Large Hadron Collider (LHC) and the discovery of a SM-like Higgs with a mass $m_h \approx 126$ GeV, seem to point towards a SUSY spectrum where at least part of the particle content have masses in the TeV range. Despite the increasing constraints on the sparticle masses and composition, a neutralino saturating the DM relic density and of almost pure Higgsino composition is able to evade current direct detection bounds due to its suppress coupling to the Higgs and Z-gauge boson [4]. Moreover, if the theory is to remain natural one expects the supersymmetric Higgs parameter $\mu \approx O(100)$ GeV, making a neutralino LSP with almost pure Higgsino composition $m_{\tilde{\chi}_1^0} \approx \mu$ a good candidate for DM. Studies have shown that pure thermal Higgsino DM is under-abundant for masses below 1 TeV [5]. This tension served as motivation for non-thermal ways of generating the right amount of Higgsino relic density [6], [7].

In this work we propose an alternative non-thermal way of generating the right amount of Higgsino DM via late decays of sneutrinos. As has been well established by now, neutrinos are massive. A convenient manner of obtaining neutrino masses is through the addition of right-handed neutrinos to the SM content, which by means of heavy Majorana masses leads to the type I see-saw mechanism of neutrino mass generation. When this extra content in the SM is supersymmetrized, we find that for the lightest sneutrino masses $m_{\tilde{\nu}_e} \gtrsim \mu$ and small Yukawa couplings $Y_N$ to the Higgsino-chargino sector, late decays of sneutrinos either directly to the lightest neutralino $\tilde{\chi}_1^0$ or cascading to it via decays to $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ can be efficient enough in generating the right amount of Higgsino DM relic density when sneutrinos are produced in the early Universe via decays of heavier SUSY particles (freeze-in scenario [8]).

The paper is organized as follows. In sec. II we briefly review the status of Higgsino DM in the MSSM and the current constraints on the parameter space. We move on in sec. III to describe the additon of the right-handed neutrino sector and how the decays of the lightest sneutrino can be efficient in generating the Higgsino DM relic density in a non-thermal way via late decays. Finally, our conclusions are given in sec. IV.

II. HIGGSINO DARK MATTER

The paradigm of Higgsino dark matter is well motivated both from arguments based on naturalness of the EW scale as well as, on a more practical sense, from current collider constrains in SUSY searches at the LHC. In the large $\tan \beta \gg 1$ limit, necessary to obtain the maximum value for the tree-level Higgs mass in the MSSM, $m_{h,\text{tree}} \approx m_Z$, the usual measure of tuning [9], $\Delta = \max (|d \log v^2/d \log \xi_i| \approx \mu^2 + m_{H_u}^2$, where $\xi_i$ are the relevant parameters of the MSSM, implies that $\mu \lesssim \text{few } O(100)$ GeV for a natural theory. Moreover, the recent discovery at the LHC of a SM-like Higgs with mass $m_h \approx 126$ GeV implies in the MSSM that large...
radiative corrections are necessary to raise the tree-level Higgs mass. In these finite-loop corrections only third generation sparticles are relevant due to their coupling to the Higgs. Thus, an effective SUSY spectrum with only third generation sparticles, a Higgsino sector and all other sparticles decoupled becomes a natural option. On the other hand, the latest SUSY searches at the LHC [10] [11], as well as flavour constraints from B-factories [12], also highly constrain first and second generation sparticles as well as gluinos, pointing towards a natural spectrum. We’d also like to point out that in the split versions of SUSY [13], [14] where a natural EW theory is no longer a requirement, a light Higgsino sector can arise due to chiral symmetry protection of the fermion masses, making Higgsino DM studies relevant for this case as well.

We concentrate on what are known as ”Higgsino-world” scenarios [15], [16] in which squarks and sleptons of the MSSM have masses in the multi-TeV range, while $\mu$ is sub-TeV. Thus, we consider the range of masses: $|\mu| \ll m_{\tilde{g}}$, $m_{\tilde{W}} \ll m_{\tilde{q}}$, $m_{\tilde{f}}$, with light Higgsino-like charginos $\tilde{\chi}_1^\pm$ and two light Higgsino-like neutralinos $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$. Though in order to get $m_h \approx 126$ GeV via a large trilinear $A_t$ the lightest stop could be sub-TeV, we assume for simplicity that the lightest stop has a mass above the TeV range. This kind of SUSY scenario has been thoroughly studied and is well-known that in the case of thermal production it leads to a very low relic density of neutralinos, in disagreement with the latest Planck results (at the 3σ level) [17]: 0.1118 $< \Omega h^2 < 0.128$. This is a consequence of the sizeable couplings involved and it implies that thermal Higgsino DM is under-abundant for $\mu \lesssim 1$ TeV [5]. Therefore, non-thermal ways of generating the correct amount of relic density have been proposed such as moduli field remnant from string theory decaying into a Higgsino-like neutralino LSP [6] or, in the midst of solving the strong-CP problem, a Peecce-Quinn axino annihilating to Higgsinos which provides a Higgsino-dominated or axion dominated DM relic density (two species of DM) depending on the which type of annihilation dominates [7]. In order to simplify our analysis, we take $m_{\tilde{W}} \sim m_{\tilde{g}}$, $m_{\tilde{f}}$, decoupling the Wino from our effective theory. It turns out that in this kind of SUSY spectrum, spin independent and spin dependent direct detection constraints can be greatly ameliorated for an almost pure Higgsino DM due to its reduced coupling to the Higgs and the Z-gauge boson [4]. Indirect detection constraints from gamma ray observations at Fermi [18] exclude non-thermal Higgsino-like DM for values of $|\mu| \lesssim 250$ GeV. Therefore an acceptable region of MSSM parameter space which satisfies all relevant DM constraints is a non-thermal mostly Higgsino DM with $|\mu| \gtrsim 250$ GeV, $M_1 \gg |\mu|$ and $\tan \beta \gg 1$, with the last constraint coming from the Higgs’ mass requirements.

### III. LATE DECAYS OF SNEUTRINOS

To generate the right amount of relic Higgsino-like DM with $m_{DM} \sim \mu \approx \mathcal{O}(100)$ GeV we need to resort to non-thermal ways. It has been well established by many experiments that at least some of the neutrinos are massive and that the different flavours oscillate in vacuum and matter. A simple way to generate neutrino masses in the SM is by adding at least 2 right-handed neutrino fields to the SM particle content. Given that this right-handed neutrinos are singlets under the SM gauge groups, a Majorana mass can be introduced for each of them which in conjunction with a Yukawa interaction involving the left-handed neutrino and the Higgs can be used in the well-known type I see-saw mechanism to generate small neutrino masses. Even in supersymmetric models we need to be able to account for massive neutrinos. In principle we could extend the MSSM to include 2 right-handed neutrino superfields and explain the solar and atmospheric neutrino mass differences. This however fixes the new Yukawa interactions between left-handed neutrinos and right-handed neutrinos and in practice does not allow late decays for the lightest sneutrino. Thus we introduce 3 right-handed neutrino superfields $N_i$ with $i=a,b,c$ to the MSSM spectrum from which we can explain the solar and atmospheric neutrino mass differences by means of two of these. The superpotential then takes the form $W = W_{MSSM} + M_{N_a} N_a \bar{N}_a + y_{N_i} L_i H_u N_i$, where $W_{MSSM}$ is the MSSM superpotential, $M_{N_i}$ the Majorana masses and $y_{N_i}$ the new Yukawa couplings. Two of these Yukawa couplings are fixed by the atmospheric and solar mass differences and we use the third Yukawa interaction to generate the late out of equilibrium decays of the corresponding sneutrino. We similarly introduce soft-breaking masses, bi-linear and tri-linear interactions for the scalar neutrinos, $\Delta L_{\text{soft}} = -m_{\tilde{N}_i}^2 |N_i|^2 + ((b_{N_i}/2)M_{N_i}N_i^2 - A_{N_i} \tilde{L}_i H_u \tilde{N}_i + h.c.)$, where we assume all couplings to be real. Since we are interested in the DM picture, we decoupled from our low energy effective theory the sneutrinos corresponding to the solution of the solar and atmospheric mass differences (with indices $i = a, b$) by taking $m_{\tilde{N}_{a,b}}^2 \sim m_{\tilde{N}_{a,b}}^2 \gg m_{\tilde{N}_c}^2$. Similarly, we assume for the Majorana masses $M_{N_a} \sim M_{N_b} \gg M_{N_c}$ decoupling the corresponding mostly right-handed neutrinos as well. Therefore, in effect we concentrate in a single superfield $N_c$ and in particular in its corresponding complex scalar component. Assuming CP-conservation in the sneutrino sector we can decouple the chiral sneutrino fields as $\tilde{\nu}_L = (\tilde{\nu}_{L,1} + i \tilde{\nu}_{L,2})/\sqrt{2}$ and $\tilde{N} = (N_1 + i N_2)/\sqrt{2}$, from now on we understand $N = N_c$, and $L = L_c$ is the corresponding left-handed lepton flavour. Then the sneutrino mass matrix reduces in the basis $(\tilde{\nu}_{L,1}, \tilde{N}_1, \tilde{\nu}_{L,2}, \tilde{N}_2)$ to a block diagonal form,

\[
\begin{pmatrix}
m_{\tilde{N}_L}^2 & m_{\tilde{N}_{RL}} + v \sin \beta Y_{N} M_N & 0 & 0 \\
m_{\tilde{N}_{RL}} - b_{N} M_N & m_{\tilde{N}_{RR}} + v \sin \beta Y_{N} M_N & 0 & 0 \\
0 & 0 & m_{\tilde{N}_L}^2 & m_{\tilde{N}_{RL}} - v \sin \beta Y_{N} M_N \\
0 & 0 & m_{\tilde{N}_{RL}} - v \sin \beta Y_{N} M_N & m_{\tilde{N}_{RR}} + b_{N} M_N \\
\end{pmatrix}
\]
where \( m_{\chi}^2 = m_{\tilde{L} L}^2 + y_N \sin^2 \beta + \left( m_{\tilde{W}}^2 / 2 \right) \cos^2 \beta \), \( m_{\tilde{R} R} = M_N^2 + m_{\tilde{N} i}^2 + y_N \sin^2 \beta \) and \( m_{\tilde{R} L} = -\mu Y_N \cos \beta + v \sin \beta \). Denoting the mass eigenstates as \( \tilde{\nu}_i \) and \( \tilde{N}_i \) with \( i = 1, 2 \), we have that \( \tilde{\nu}_i = \tilde{\nu}_i + \tilde{N}_i \sin \theta_i \) and \( \tilde{N}_i = \tilde{\nu}_i \cos \theta_i \), where,

\[
\tan 2\theta_i = \frac{2(m_{\tilde{R} L} - v \sin \beta M_N Y_N)}{m_{\tilde{L} L} - (m_{\tilde{R} R} + b_M M_N)}. \tag{3.1}
\]

At this stage we discuss the requirements for the non-thermal Higgsino-like DM generation from sneutrino late decays. Calling the lightest sneutrino mass-eigenstate \( \tilde{\nu}_0 \), we choose it such that it corresponds to the CP-even sneutrino sector \( \tilde{\nu}_1 \) in Fig. 1, the dependence of the di-tau threshold. We calculated the 3-body decay into the second lightest neutralino \( \tilde{\chi}_1^0 \) related to the sneutrino sector \( \nu \) in this scenario are governed effectively by the Hubble expansion rate evaluated at the \( \tilde{\nu}_0 \) decay time. Assuming that no significant entropy is generated between the \( \tilde{\nu}_0 \) decays till nowadays and that the Universe is radiation dominated at the decay time epoch, the condition takes the form,

\[
0.1 \frac{\rho_{\,h}}{m_{\chi}^2} \frac{s(x_d)}{s_0} \frac{|\sigma v(x_d)|}{|\sigma v^{\chi \chi 
 \rightarrow \chi \chi _{\text{thermal}}}|} \frac{m_{\chi}^2}{M_{\chi}^2} < 1.67 \sqrt{g_*} \frac{1}{M_{\chi}^2 x_d^2}. \tag{3.3}
\]

where \( x_d = m_{\chi}^2 / T_d \) is related to the temperature at which the decays happen, \( \langle \sigma v(x_d) \rangle \frac{\chi \chi \rightarrow \chi \chi _{\text{thermal}} \rightarrow \text{SM} \rangle \) is the annihilation rate of the LSP into SM particles, \( \rho_{\,h} = 8.06 \times 10^{-47} h^2 \text{GeV}^4 \) is the critical density of the Universe, \( s_0 = 2.22 \times 10^{-38} \text{GeV}^3 \) is the entropy of the Universe today, \( g_* \) are the active degrees of freedom and \( s(x_d) = (2\pi^2 / 45) g_* (m_{\chi}^3 / x_d) \) the entropy evaluated at the time of decay. Taking the inequality Eq. (3.3) and solving for \( x_d \), we find that

\[
x_d \gtrsim \frac{2.6 \times 10^{-2} \sqrt{g_* M_{\chi} (r_{\,h}) \langle \sigma v(x_d) \rangle |\chi \chi \rightarrow \chi \chi _{\text{thermal}} \rightarrow \text{SM}|}}{s_0}. \tag{3.4}
\]

Typical thermal relic densities for a Higgsino-like LSP are of the order of \( \Omega_{\chi}^2 \sim 10^{-2} \), see for example [5], [19]. Since the Higgsino annihilations are s-wave dominated, we can get an estimate of the corresponding thermally averaged annihilation cross-sections \( \langle \sigma v(x_d) \rangle \frac{\chi \chi \rightarrow \text{SM} \rangle \sim 2 \times 10^{-8} \text{GeV}^{-2} \) and therefore on \( x_d \) via Eq. (3.4), \( x_d \gtrsim 230 \). We see that for values

[1] We could have similarly chosen it to correspond to the CP-odd sector.
of $m_\tilde{\nu}_0 \sim \mu \gtrsim 300$ GeV, this corresponds to a decay temperature $T_D \lesssim 1$ GeV or similarly to a decay time $t_d \gtrsim 10^{-6}$ s, safely below Big Bang Nucleosynthesis (BBN) times.

The demand of a sufficiently late decay of sneutrinos $\tilde{\nu}_0$ implies in particular that for $m_{\nu_0} \gtrsim \mu$, $Y_N \lesssim 10^{-10}$ as can be seen from Fig. 1, and that $|\sin \theta_1| \lesssim 2 \times 10^{-5}$ if we choose the corresponding charged lepton to be the electron ($Y_L = Y_E \approx 10^{-6}$). Thus, the associated light neutrino $\nu_H$ is basically massless, while the heavy neutrino $\nu_H$ is mostly right-handed with a mass $m_{\nu_H} = M_N$, which we take to be $m_{\nu_H} \gtrsim m_{\tilde{\nu}_1}$. Notice that decays of $\nu_H \rightarrow \tilde{\nu}_0 h$ are suppressed by $\Gamma_{\nu_H \rightarrow \tilde{\nu}_0 h} \propto \sin^2 \theta_1 Y_N^2$, making their contribution negligible. The sneutrino $\tilde{\nu}_0$ is also highly right-handed and interacts minimally with early Universe plasma. Self-annihilations and possible co-annihilations cross-sections are all too small to reproduce the correct $\chi^0_1$-relic density, suppressed by either $\sin^2 \theta_1$, $\sin^4 \theta_1$, $Y_N^2$ or combinations of these, see Ref.[20]. Entropy generation in the decays is minimal since the Universe is radiation dominated and thus decays are not relevant in reducing the DM relic density. We conclude that the demand of sufficient late decays for the sneutrinos and the appropriate value of the $\chi^0_1$ relic density moves us to consider a model where the sneutrinos $\tilde{\nu}_0$ density is generated via decays of heavier SUSY particles in what is known as an example of the freeze-in mechanism [8], [21], [22]. Since we want small mixing angles $\sin \theta_1 \ll 1$, we assume that $m_{\tilde{\nu}_1L} \gg m_{\tilde{\nu}_1R}$ and $m_{\tilde{\nu}_1L} \ll m_{\tilde{\nu}_1L}$, making the mostly left-handed sneutrinos heavy $m_{\tilde{\nu}_1L} \gg m_{\nu_0}$. Furthermore we appropriately choose $b_N$ such that $m_{\tilde{\nu}_1} > m_{\tilde{\nu}_1} \approx m_{\tilde{\nu}_0}$. The relevant decays to produce $\tilde{\nu}_0$-sneutrinos are $\tilde{\nu}_1 \rightarrow \tilde{\nu}_0 h$ and $\tilde{l}^\pm \rightarrow \tilde{\nu}_0 W^\pm$ since these are the only sparticles with which $\tilde{\nu}_0$ couples to. We can solve for the "would-be" relic density of $\tilde{\nu}_0$ integrating the Boltzmann equation from the re-heating epoch ($T_R$) till the time when the reaction rate decouples ($T_F$) to find,

$$\Omega_{\tilde{\nu}_0} = \sum_{i=\tilde{\nu}_1, \tilde{l}^\pm} \frac{m_{\nu_0}}{\rho_c} \frac{s_0}{s(x_F,i)} \frac{1}{H(m_i)} \int_{x_{i,R}}^{x_{F,i}} dx_i x_i^3 C[x_i]$$ \hspace{2cm} (3.5)

where $H(m) = 1.67 \sqrt{g_0} m^2 / M_{Pl}$ and $C[x]$ is the collision term, which for decays takes the form,

$$C[x] = \int e^{-\frac{E_i}{K_i}} \frac{d^3 p_i}{(2\pi)^3}$$ \hspace{2cm} (3.6)

with $\Gamma_i = \Gamma_i^{CM} / \gamma_i$, the decay rate related to the center of mass (CM) decay rate via the time dilation factor $\gamma_i = E_i / m_i$. In the case at hand, where $\Gamma_i^{CM}$ is a constant, we can readily do the integral Eq. (3.6) and obtain $C[x_i] = m_i^3 \Gamma_i^{CM} K_1(x_i) / (2\pi^2 x_i)$, where $K_1(x_i)$ is the modified Bessel function of the second kind. For $T_R \gg m_{\tilde{\nu}_1}, m_{\tilde{l}^\pm}$, we get an approximate solution for the integral,

$$\int_{x_{R,i}}^{x_{F,i}} dx_i x_i^4 K_1(x_i) x_i \approx \frac{3}{2} \pi - \frac{m_i^3}{2T_F^3}$$ \hspace{2cm} (3.7)

which shows that the contributions from decays is insensitive to the re-heating temperature as long as it is large enough. The expression for the $\tilde{\nu}_1 \rightarrow \tilde{\nu}_0 h$ decay rate in the CM frame is,

$$\Gamma_{\tilde{\nu}_1 \rightarrow \tilde{\nu}_0 h} = \frac{|C_{\tilde{\nu}_1 \tilde{\nu}_0 h}|^2}{16\pi m_{\tilde{\nu}_1}^4} \left[ 1 - \frac{m_h^2 + m_{\tilde{\nu}_0}^2 + (m_{\tilde{\nu}_0}^2 - m_h^2)^2}{m_{\tilde{\nu}_1}^2} \right]$$ \hspace{2cm} (3.8)

where $C_{\tilde{\nu}_1 \tilde{\nu}_0 h} \approx A_N (\cos^2 \theta_1 - \sin^2 \theta_1)$ is the $\tilde{\nu}_1 \tilde{\nu}_0 h$-coupling. Plugging this expression into Eqs. (3.5, 3.6) and calculating the relic density of $\chi^0_1$ we find that in order to get a DM relic density $\Omega_{\chi^0_1 h^2} \sim 0.1$, we need $A_N \sim 10^{-9}$, independently of tan $\beta$ and only mildly dependent on $\tilde{\nu}_1$, $m_{\tilde{\nu}_0}$ and $m_{\tilde{\nu}_1}$. This at the same time implies that $\sin \theta_1 \lesssim 10^{-12}$, making the contribution from the decay rate $\Gamma_{\tilde{l}^\pm \rightarrow \tilde{\nu}_0 W^\pm} \propto \sin^2 \theta_1$ negligible. Notice that one naturally obtains a small value of $A_N$ with the usual SUSY-breaking universality conditions, $A_N \propto Y_N$, due to the late decay condition, $Y_N \sim O(10^{-10})$. We show in Fig. 2 a typical region of parameter space where the correct DM relic density is obtained.

![FIG. 2. Region of $m_{\tilde{\nu}_1} [GeV]$ vs $A_N [GeV]$ where 0.1118 < $\Omega_{\chi^0_1 h^2} < 0.128$ in agreement with Planck, with $m_{\chi^0_1} = 300$ GeV and $m_{\tilde{\nu}_0} = 500$ GeV.

IV. CONCLUSION

Motivated by the type I see-saw mechanism for the generation of neutrino masses, we have shown an alternative non-thermal way of generating Higgsino-like DM in a trivial extension of MSSM via late decays of highly
sterile, mostly right-handed sneutrino $\tilde{\nu}_0$. Due to the smallness of the $\tilde{\nu}_0$-interactions, in order not to overclose the energy density of our Universe, we are forced to consider that the sneutrino $\tilde{\nu}_0$ is never in thermal equilibrium with the early Universe plasma, and thus its number density is produced via the freeze-in process in the form of decays of heavier SUSY particles. This condition turns out to be natural for $A_N \propto Y_N$, given that $Y_N \lesssim 10^{-10}$, allowing a correct Higgsino-like DM relic density to be obtained in accordance with the latest Planck measurements.

**ACKNOWLEDGEMENTS**

We thank Michael A. Schmidt, Timothy Trott and Carlos E. M. Wagner for helpful discussions and comments. ADM is supported by the Australian Research Council.

[1] D. Akerib et al. (LUX Collaboration), Phys.Rev.Lett. **112**, 091303 (2014), 1310.8214.
[2] E. Aprile et al. (XENON100 Collaboration), Phys.Rev.Lett. **109**, 181301 (2012), 1207.5988.
[3] M. Aartsen et al. (IceCube collaboration), Phys.Rev.Lett. **110**, 131302 (2013), 1212.4097.
[4] C. Cheung, L. J. Hall, D. Pinner, and J. T. Ruderman, JHEP **1305**, 100 (2013), 1211.4873.
[5] M. Cirelli, N. Fornengo, and A. Strumia, Nucl.Phys. **B753**, 178 (2006), hep-ph/0512090.
[6] B. S. Acharya, G. Kane, S. Watson, and P. Kumar, Phys.Rev. **D80**, 083529 (2009), 0908.2430.
[7] H. Baer, A. Lessa, S. Rajagopalan, and W. Sreethawong, JCAP **1106**, 031 (2011), 1103.5413.
[8] L. J. Hall, K. Jedamzik, J. March-Russell, and S. M. West, JHEP **1003**, 080 (2010), 0911.1120.
[9] R. Barbieri and G. Giudice, Nucl.Phys. **B306**, 63 (1988).
[10] G. Aad et al. (ATLAS Collaboration), JHEP **1406**, 035 (2014), 1404.2500.
[11] S. Chatrchyan et al. (CMS Collaboration), JHEP **1401**, 163 (2014), 1311.6736.
[12] Y. Amhis et al. (Heavy Flavor Averaging Group) (2012), 1207.1158.
[13] N. Arkani-Hamed and S. Dimopoulos, JHEP **0506**, 073 (2005), hep-th/0405159.
[14] A. Arvanitaki, N. Craig, S. Dimopoulos, and G. Villadoro, JHEP **1302**, 126 (2013), 1210.0555.
[15] G. L. Kane (1998).
[16] H. Baer, V. Barger, and P. Huang, JHEP **1111**, 031 (2011), 1107.5581.
[17] P. Ade et al. (Planck Collaboration), Astron.Astrophys. (2014), 1303.5076.
[18] M. Ackermann et al. (Fermi-LAT collaboration), Phys.Rev.Lett. **107**, 241302 (2011), 1108.3546.
[19] M. Cirelli, A. Strumia, and M. Tamburini, Nucl.Phys. **B787**, 152 (2007), 0706.4071.
[20] S. Gopalakrishna, A. de Gouvea, and W. Porod, JCAP **0605**, 005 (2006), hep-ph/0602027.
[21] K. Petraki and A. Kusenko, Phys.Rev. **D77**, 065014 (2008), 0711.4646.
[22] A. D. Medina and E. Ponton, JHEP **1109**, 016 (2011), 1104.4124.