Two Puzzles, One Solution: Neutrino Mass and Secluded Dark Matter

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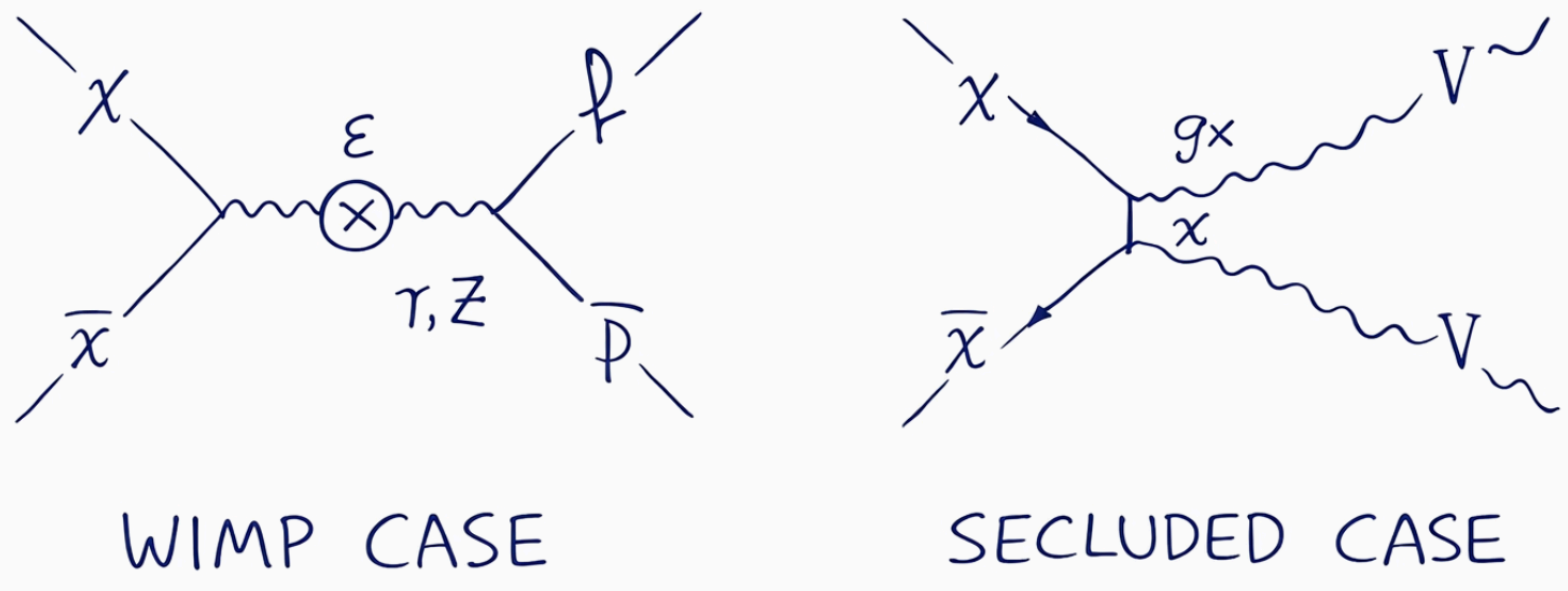
- M. DI Mauro. arXiv:2511.19622 Two Puzzles, One Solution: Neutrino Mass and Secluded Dark Matter. Accepted in PRD
- M.Di Mauro, Yanhan Wang. arXiv:2510.23771, WIMP Shadows: Phenomenology of Secluded Dark Matter in Three Minimal BSM Scenarios. Accepted in PRD

Istituto Nazionale di Fisica Nucleare

https://theconversation.com/why-do-astronomers-believe-in-dark-matter-122864

Moriond2026, EW, La Thuille March 20th

Secluded dark matter models



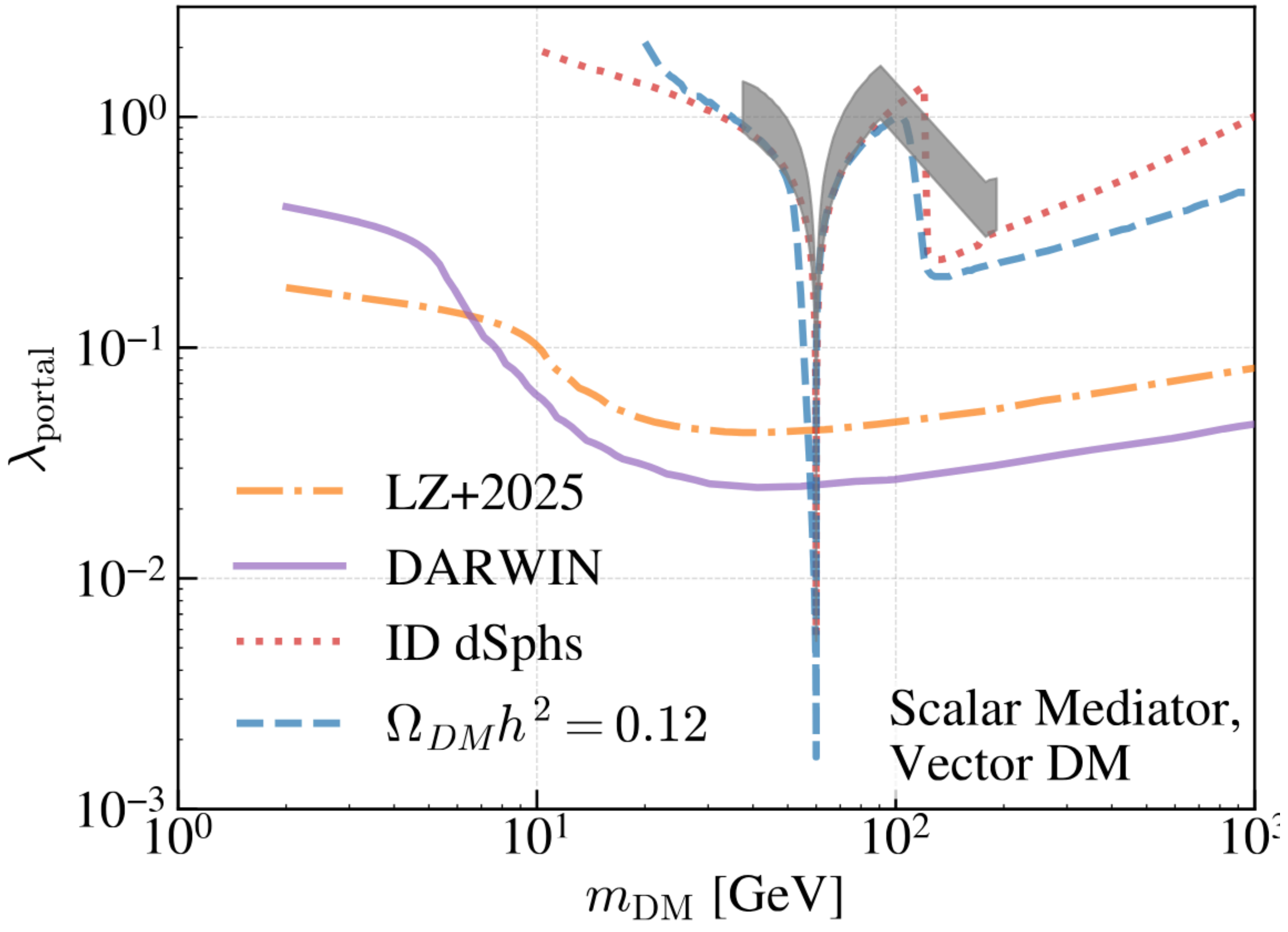
WIMP CASE: $\varepsilon \approx O(1)$

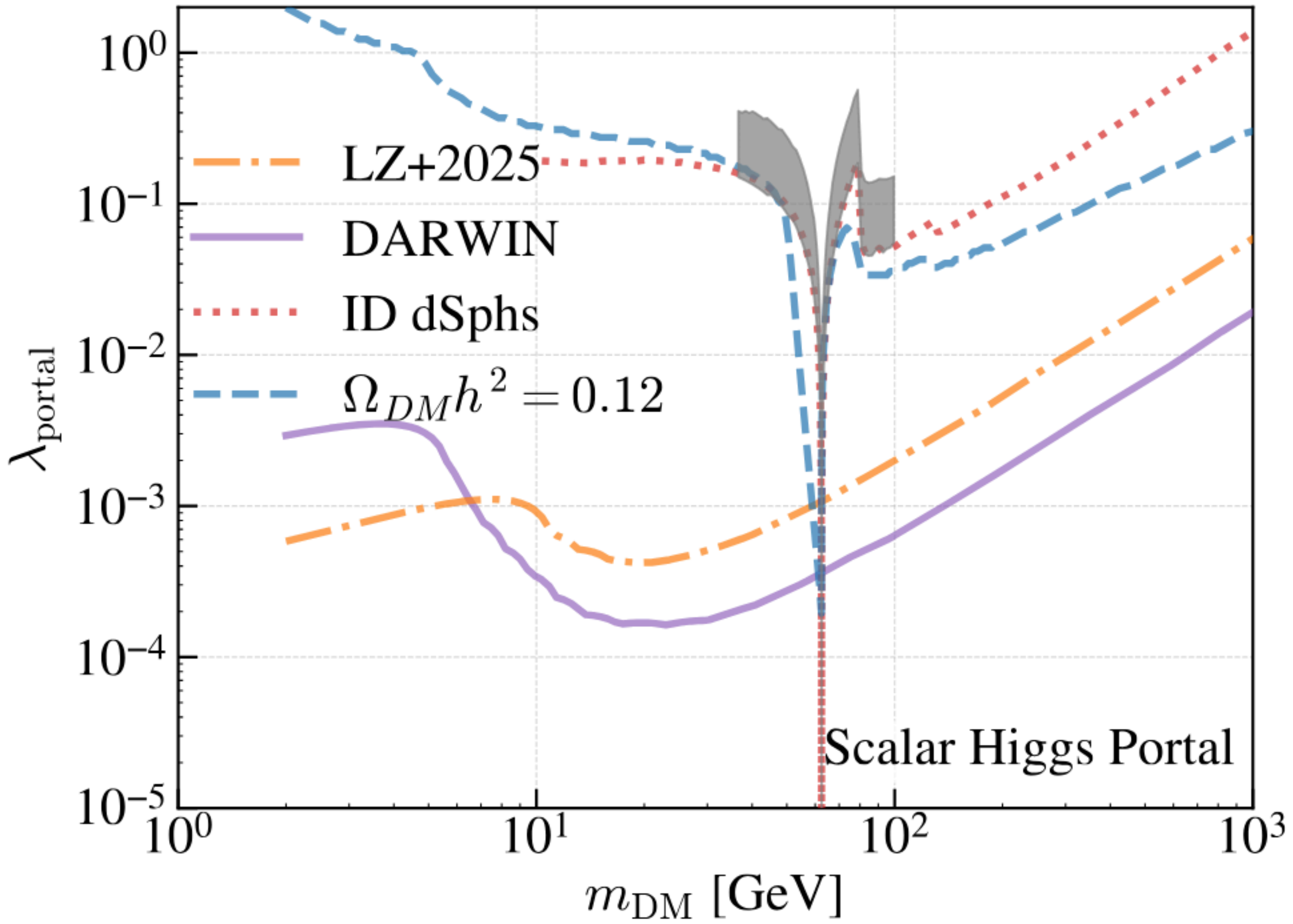
SECLUDED CASE: $\varepsilon \ll 1 \quad g_\chi \sim O(0,1)$

$m_\chi > m_V$

See also Giorgio Arcadi's talk

Kong, Di Mauro *Phys.Rev.D* 113 (2026) 4, 043031





Pospelov, Ritz, Voloshin
Phys.Lett.B 662 (2008) 53-61

Three extensions of the standard model

Kinetic mixing

$$\mathcal{L} \supset \bar{\chi}(i\not{D} - m_\chi)\chi - \frac{1}{4}F'_{\mu\nu}F'^{\mu\nu} + \frac{1}{2}m^2_{Z'}Z'_\mu Z'^\mu + \frac{\epsilon}{2}F'_{\mu\nu}B^{\mu\nu}$$

Simplified model

$$\mathcal{L} \supset \partial_\mu K \partial^\mu K - m_K^2|K|^2 + \frac{1}{2}\partial_\mu S \partial^\mu S - \frac{1}{2}m_S^2 S^2 + m_K g_K K^3 + m_S \lambda_{KS} S^2 K + \lambda_{KS}^2 S^2 K^2 + \frac{\beta}{\sqrt{2}}\frac{m_f}{v_h}\bar{f}fK.$$

Higgs Portal

$$\mathcal{L} \supset -\frac{1}{4}F'_{\mu\nu}F'^{\mu\nu} - \frac{\epsilon}{2}F'_{\mu\nu}B^{\mu\nu} + (D_\mu R)^\dagger(D^\mu R) - V(\Phi, R) + \bar{\chi}(i\not{D} - m_\chi)\chi - \left(y_p\,\bar{\chi}\chi\,R + \text{h.c.}\right),$$

$$V(\Phi, R) = \mu_H^2\,\Phi^\dagger\Phi + \mu_R^2\,R^\dagger R + \lambda_H(\Phi^\dagger\Phi)^2 + \lambda_R(R^\dagger R)^2 + \lambda_{HR}(\Phi^\dagger\Phi)(R^\dagger R).$$

Di Mauro, Wang arXiv:2510.23771

Thermal history of the mediator

- decays and inverse decays, $\boxed{\phi \leftrightarrow \mathrm{SM\,SM}}$, with thermally averaged rate

$$\langle \Gamma_\phi \rangle (T) \equiv \Gamma_\phi \, \frac{K_1(m_\phi/T)}{K_2(m_\phi/T)} \,,$$

where Γ_ϕ is the ϕ decay width at rest;

- annihilations and inverse annihilations into SM particles, $\boxed{\phi\phi \leftrightarrow \mathrm{SM\,SM}}$ (or $\phi\phi^\dagger$ if ϕ is complex), with rate

$$\Gamma_{\mathrm{ann}}(T) \equiv n_\phi^{\mathrm{eq}}(T) \, \langle \sigma v_{\mathrm{rel}} \rangle_{\phi\phi \to \mathrm{SM\,SM}} \,,$$

where n_ϕ^{eq} is the equilibrium density of ϕ;

- elastic scattering, $\boxed{\phi f \leftrightarrow \phi f}$, with rate

$$\Gamma_{\mathrm{el}}(T) \equiv \sum_f n_f^{\mathrm{eq}}(T) \, \langle \sigma v_{\mathrm{rel}} \rangle_{\phi f \to \phi f} \,,$$

which maintains kinetic equilibrium but does not change n_ϕ.

Keeping the *number density* near equilibrium, $n_\phi \simeq n_{\phi,\mathrm{eq}}$, additionally requires *chemical* processes, such as decay/inverse decay of ϕ and annihilation into SM particles, to be faster than the Hubble rate H:

$$\Gamma_{\mathrm{chem}}(T) \;\equiv\; \langle \Gamma_\phi \rangle (T) + n_\phi^{\mathrm{eq}}(T) \, \langle \sigma v_{\mathrm{rel}} \rangle_{\phi\phi \to \mathrm{SM\,SM}} \;\gtrsim\; H(T)$$

The $2 \to 2$ *elastic* scattering of ϕ with the bath (e.g. $\phi f \leftrightarrow \phi f$) maintains *kinetic* equilibrium provided

$$\Gamma_{\mathrm{el}}(T) \;\equiv\; \sum_f n_f^{\mathrm{eq}}(T) \, \langle \sigma v_{\mathrm{rel}} \rangle_{\phi f \to \phi f} \;\gg\; H(T) \,.$$

Di Mauro, Wang arXiv:2510.23771

Thermal equilibrium of the dark-sector mediator

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle\sigma v_{\rm rel}\rangle_{\chi\bar\chi\to\phi\phi}\left(n_\chi^2 - n_{\chi,{\rm eq}}^2\frac{n_\phi^2}{n_{\phi,{\rm eq}}^2}\right) - \langle\sigma v_{\rm rel}\rangle_{\chi\bar\chi\to{\rm SM}}\left(n_\chi^2 - n_{\chi,{\rm eq}}^2\right),$$

$$\frac{dn_\phi}{dt} + 3Hn_\phi = +\langle\sigma v_{\rm rel}\rangle_{\chi\bar\chi\to\phi\phi}\left(n_\chi^2 - n_{\chi,{\rm eq}}^2\frac{n_\phi^2}{n_{\phi,{\rm eq}}^2}\right) - \langle\sigma v_{\rm rel}\rangle_{\phi\phi\to{\rm SM}}\left(n_\phi^2 - n_{\phi,{\rm eq}}^2\right) - \langle\Gamma_\phi\rangle\left(n_\phi - n_{\phi,{\rm eq}}\right).$$

↓

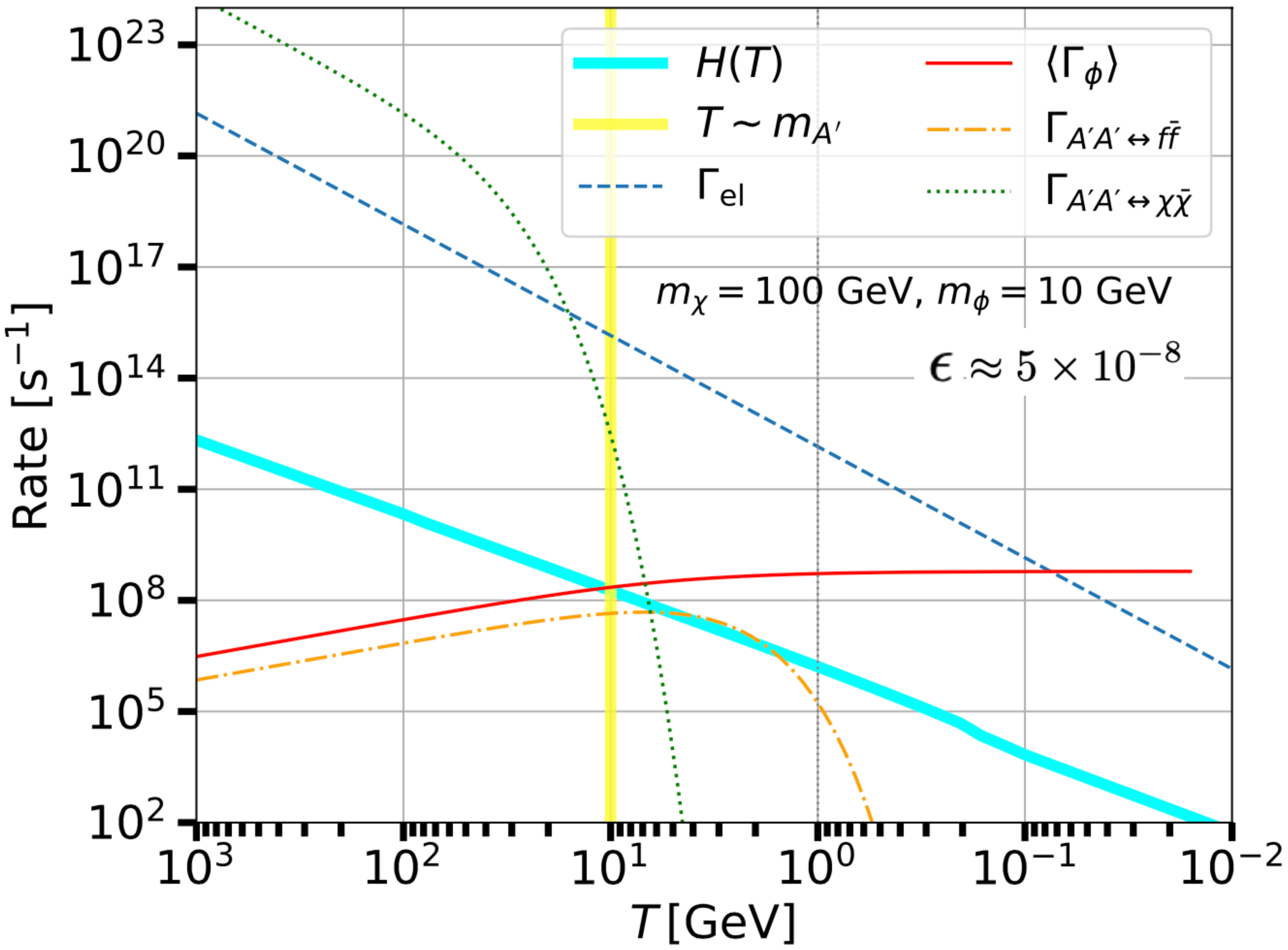
$$\Gamma_{\rm chem}(T) \equiv \langle\Gamma_\phi\rangle(T) + n_\phi^{\rm eq}(T)\,\langle\sigma v_{\rm rel}\rangle_{\phi\phi\to{\rm SM\,SM}} \gtrsim H(T) \qquad \Gamma_{\rm el}(T) \equiv \sum_f n_f^{\rm eq}(T)\,\langle\sigma v_{\rm rel}\rangle_{\phi f\to\phi f} \gg H(T)$$

↓

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle\sigma v_{\rm M\text{ø}l}\rangle_T\left(n_\chi^2 - n_{\chi,{\rm eq}}^2\right)$$

Di Mauro, Wang arXiv:2510.23771

Thermal equilibrium of the dark-sector mediator



$$\mathcal{L} \supset \bar{\chi}(i\not{D} - m_\chi)\chi - \frac{1}{4}F'_{\mu\nu}F'^{\mu\nu} + \frac{1}{2}m_{Z'}^2 Z'_\mu Z'^\mu + \frac{\epsilon}{2}F'_{\mu\nu}B^{\mu\nu}$$

$$\epsilon \gtrsim \left[\frac{1.66\sqrt{g_*}}{\frac{\alpha}{3}}\,\frac{m_{A'}}{M_{\rm Pl}}\,\frac{1}{(K_1/K_2)|_{x_{A'}=1}}\right]^{1/2} \approx$$
$$\approx 5 \times 10^{-8}\left(\frac{g_*}{80}\right)^{1/4}\left(\frac{m_{A'}}{10\ \mathrm{GeV}}\right)^{1/2}.$$

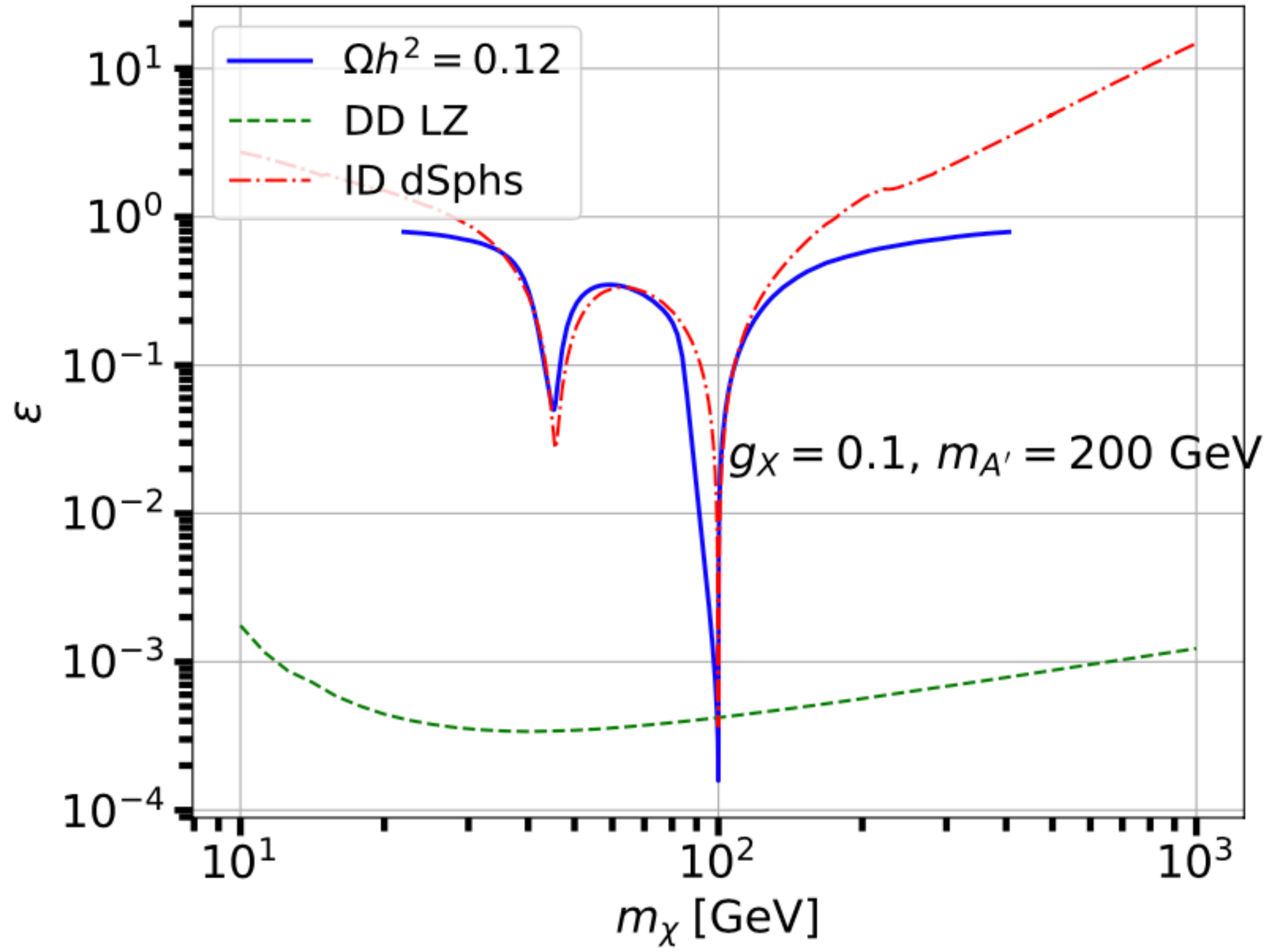
BBN constraints

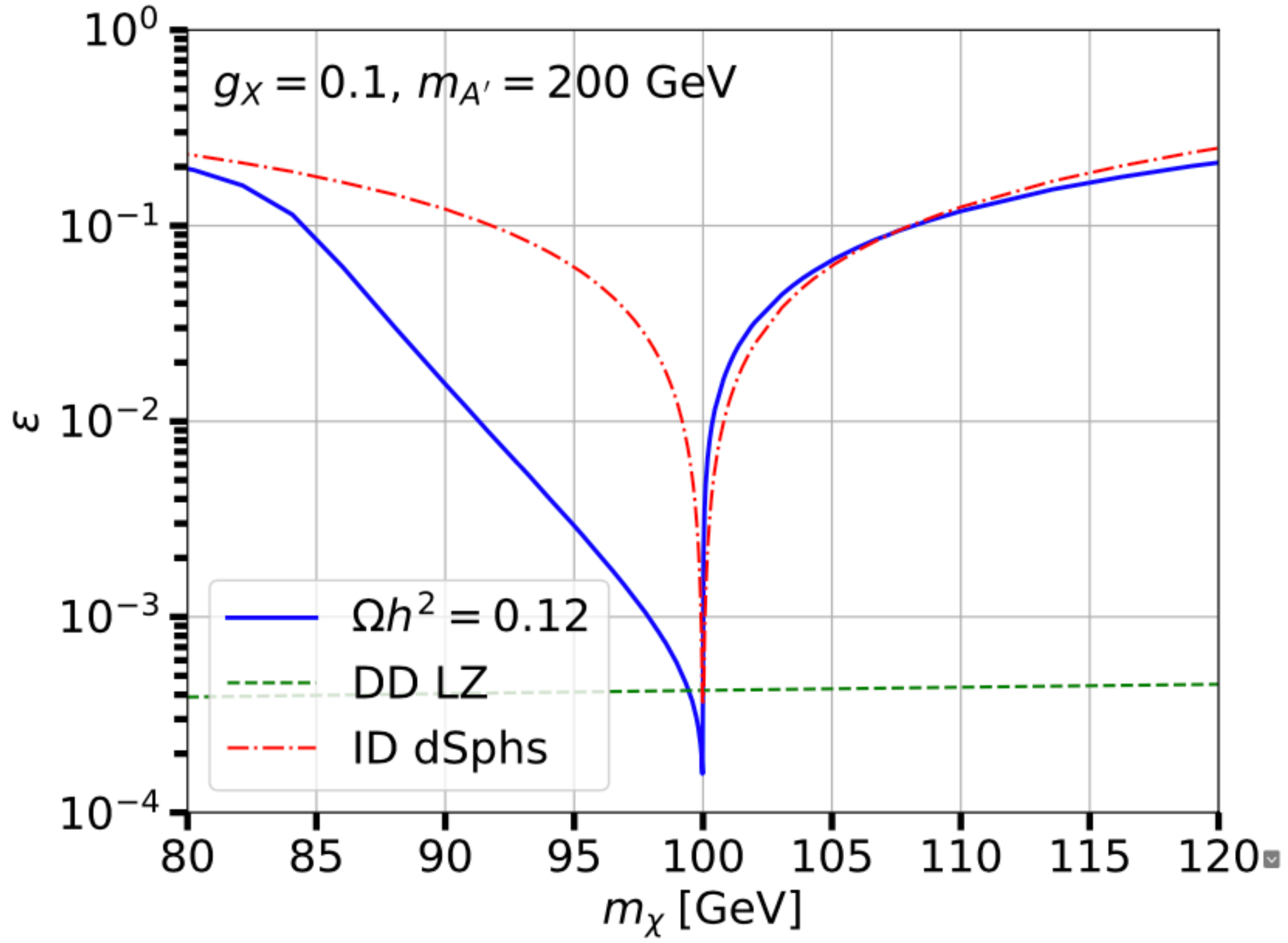
$$\tau_\phi \lesssim 1\ \mathrm{s}\ (T \gtrsim 1\ \mathrm{MeV}) \Longleftrightarrow \Gamma_\phi \gtrsim 1\ \mathrm{s}^{-1} \simeq 6.6\times10^{-25}\ \mathrm{GeV}$$

$$\epsilon \gtrsim 5.2 \times 10^{-11}\sqrt{\frac{1\ \mathrm{GeV}}{m_{A'}}}\left[\frac{1}{\sum_{f\ \mathrm{open}} N_c^f Q_f^2}\right]^{1/2}$$

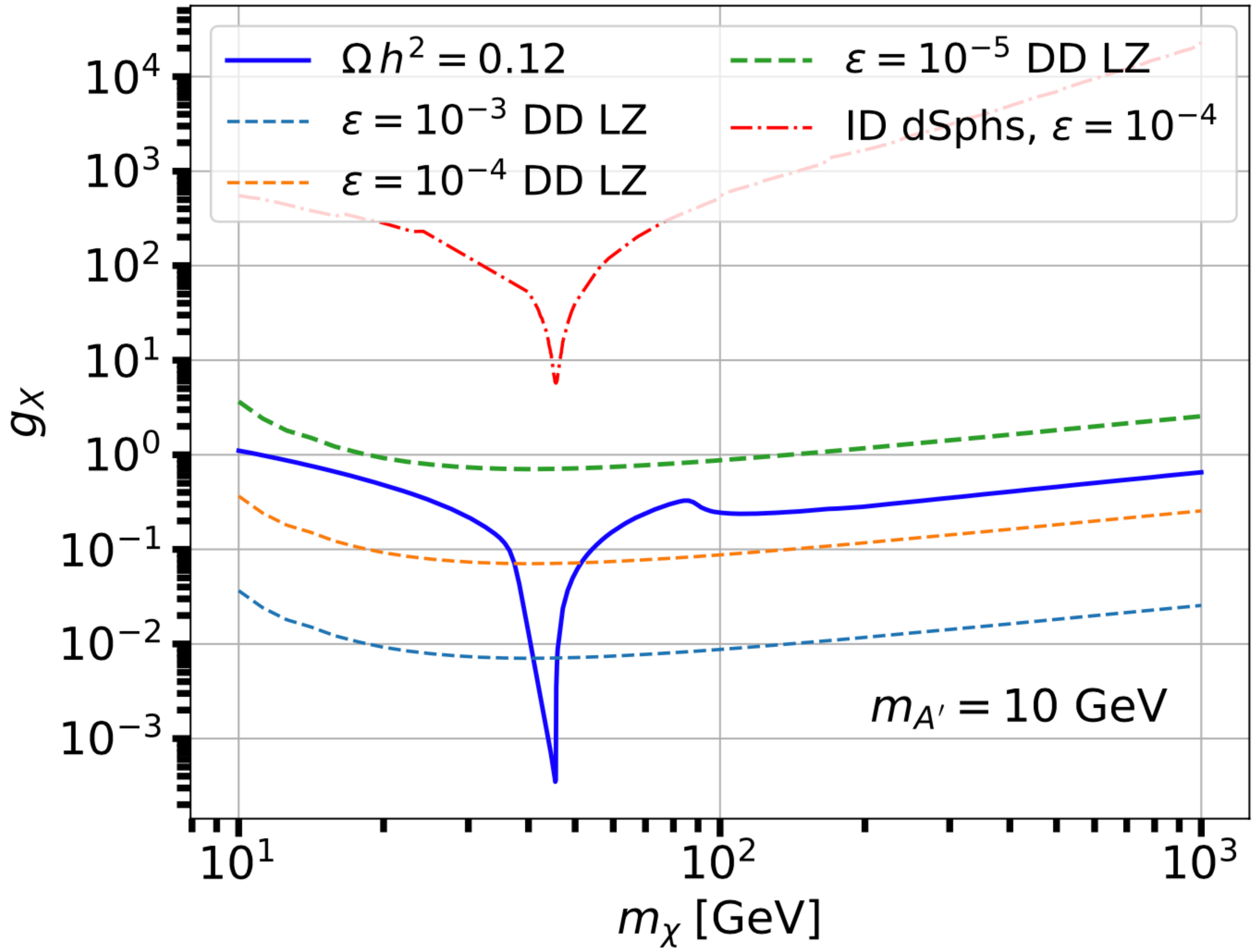
Model I (Dirac dark matter and an extra U(1)x)

WIMP CASE





SECLUDED CASE



Di Mauro, Wang arXiv:2510.23771

Connecting Neutrino mass and Dark Matter ‘invisibility’

- Is there a theoretical reason why ε<<1?
- Another quantity in the SM model is extremely small and we do not know why: The neutrino mass that is 10^{-7} times lighter than the electron mass.

$$\mathcal{L} \supset -\frac{1}{4} F'_{\mu\nu} F'^{\mu\nu} - \frac{\epsilon}{2} F'_{\mu\nu} B^{\mu\nu} + \partial_\mu R \partial^\mu R - V(\Phi, R) + \bar{\chi}(i\not{D} - m_\chi)\chi - \left(y_p \bar{\chi}\chi R + Y_\nu^{\alpha i} \bar{L}_\alpha \tilde{\Phi} N_i + \tfrac{1}{2} Y_N^{ij} R \overline{N_i^c} N_j + \text{h.c.} \right)$$

$$\begin{aligned} V(\Phi, R) &= \mu_H^2 \Phi^\dagger \Phi + \mu_R^2 R^2 + \lambda_H (\Phi^\dagger \Phi)^2 \\ &+ \lambda_R R^4 + \kappa (\Phi^\dagger \Phi) R^2 . \end{aligned}$$

$$\Phi = \begin{pmatrix} G^+ \\ \frac{v_h + h + iG^0}{\sqrt{2}} \end{pmatrix}, \qquad R = \frac{v_r + \rho}{\sqrt{2}},$$

$$\begin{pmatrix} H \\ H_p \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} h \\ \rho \end{pmatrix}$$

$$\mathcal{L}_{H,H_p-f\bar{f}} = -\sum_f \frac{m_f}{v_h} \left(\cos\alpha\, H - \sin\alpha\, H_p \right) \bar{f} f \qquad \mathcal{L}_{H,H_p-\chi\bar{\chi}} = -\frac{y_p}{\sqrt{2}} \left(\sin\alpha\, H + \cos\alpha\, H_p \right) \bar{\chi}\chi$$

$$\mathcal{L}_{H,H_p-NN} = -\sum_{I=1}^{3} \frac{M_I}{2 v_r} \left(\sin\alpha\, H + \cos\alpha\, H_p \right) \overline{N_I^c} N_I$$

M. DI Mauro. arXiv:2511.19622

See-saw mechanism

$$\mathcal{L}_Y \supset -\, Y_\nu^{\alpha i}\, \overline{L_\alpha}\, \tilde{\Phi}\, N_i - \frac{1}{2} Y_N^{ij}\, R\, \overline{N_i^c} N_j + \text{h.c.} \qquad m_D \equiv \frac{v_h}{\sqrt{2}} Y_\nu, \qquad M_N \equiv \frac{v_r}{\sqrt{2}} Y_N$$

$$M_\nu \;\simeq\; -\, m_D\, M_N^{-1}\, m_D^T$$

$$m_{\text{light}} \simeq -\frac{m_D^2}{M_N}, \qquad m_{\text{heavy}} \simeq M_N \left(1 + \frac{m_D^2}{M_N^2}\right)$$

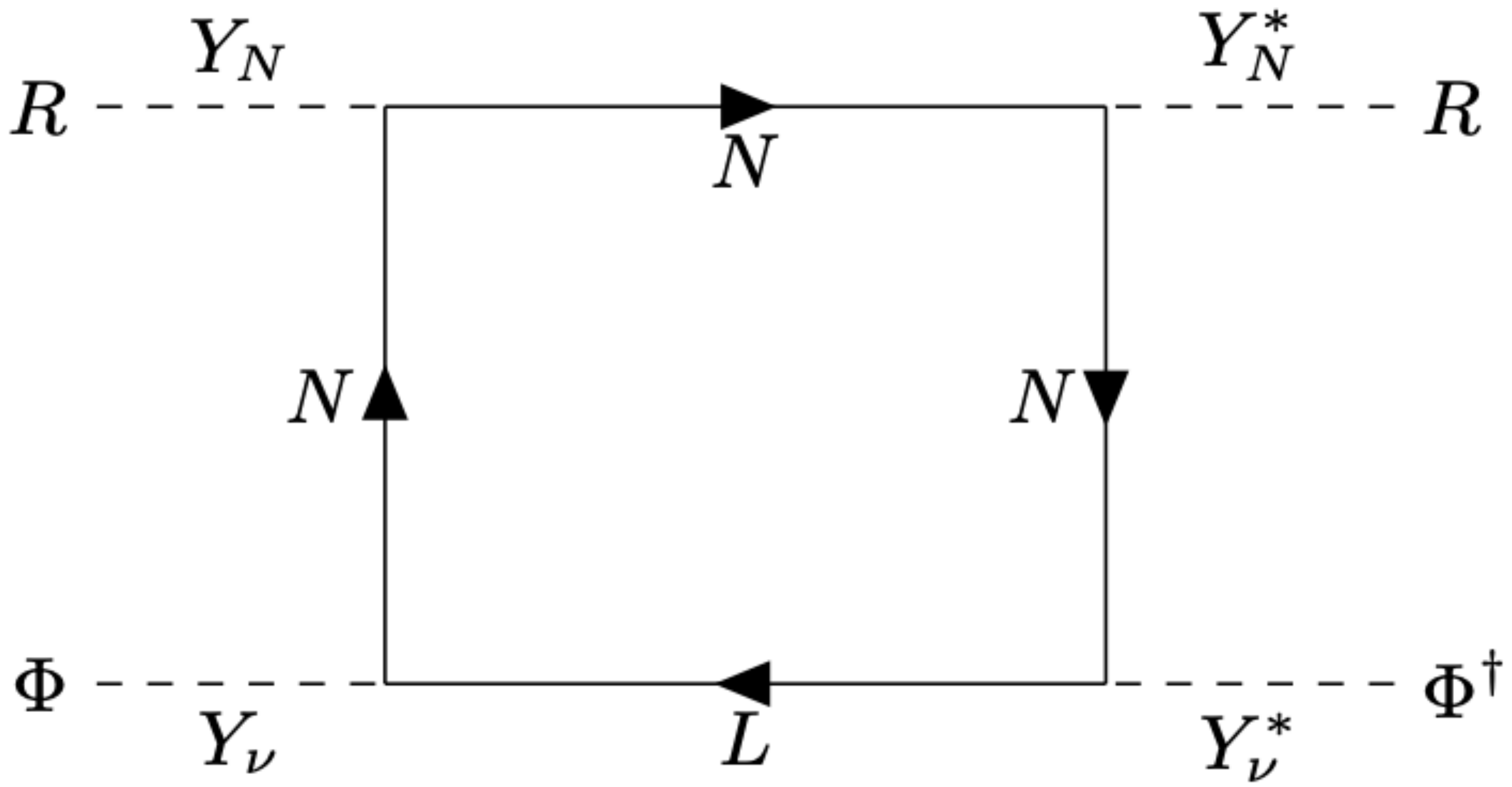
Casas–Ibarra in the Aligned case

$$m_D = i\, U_\nu \sqrt{\widehat{m}_\nu}\, O \sqrt{\widehat{M}_N}, \qquad O^T O = \mathbb{1} \qquad \widetilde{m}_I \equiv \frac{(m_D^\dagger m_D)_{II}}{M_I} = m_{\nu,I} = \frac{v_h^2}{2} \frac{|y_{\nu,I}|^2}{M_I}$$

M. DI Mauro. arXiv:2511.19622

Loop induced portal

$$\kappa(\Phi^\dagger\Phi)R^2 \qquad \kappa(\Lambda) = 0$$

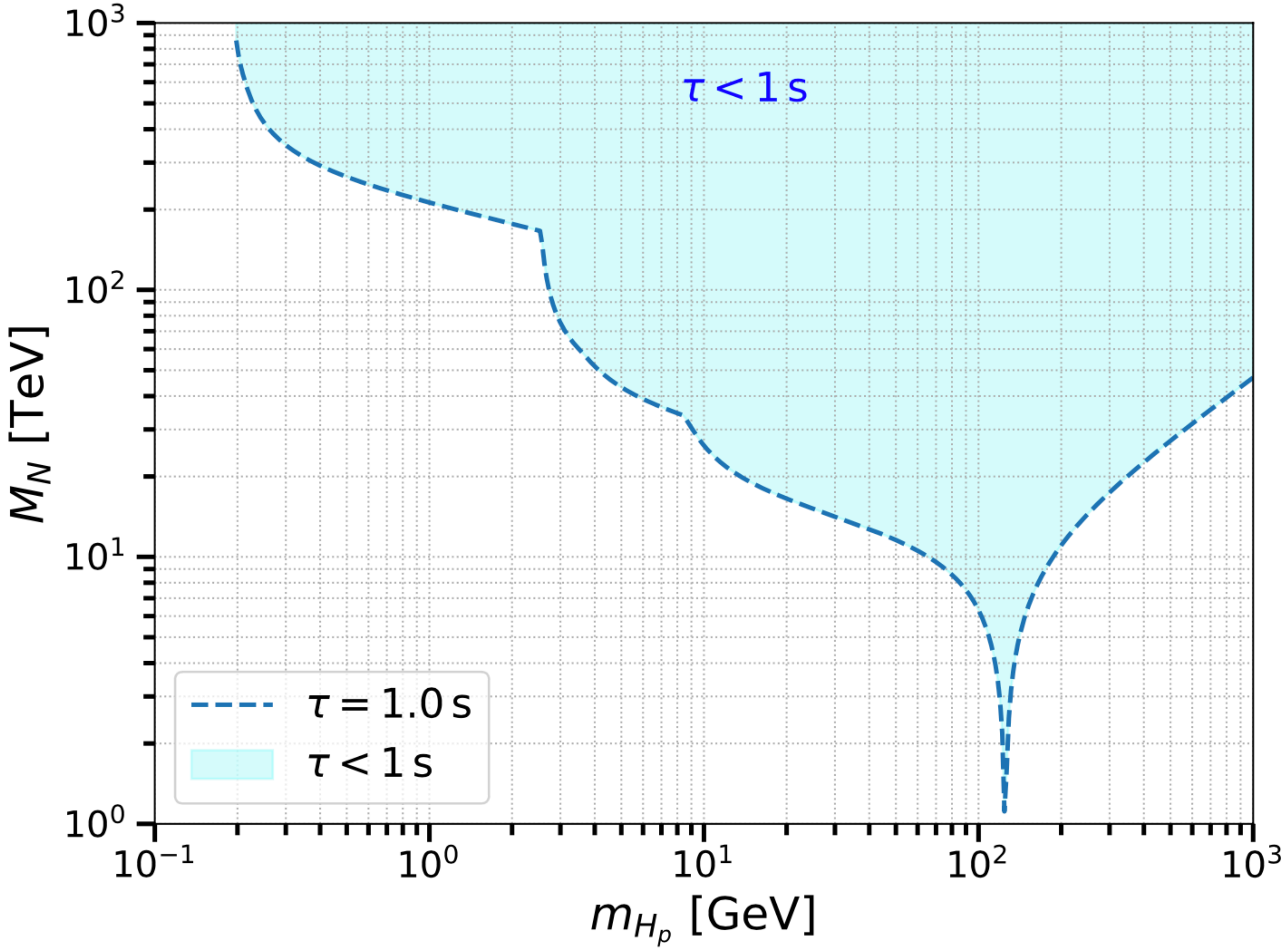


$$\kappa_{\text{loop}} = -\sum_{I=1}^{3} \frac{y_{N,I}^2\, M_{N,I}}{8\pi^2 v_h^2}\, m_{\nu,I},$$

$$\tan 2\alpha = \frac{2\,\kappa\, v_h v_r}{m_{H_p}^2 - m_H^2} = -\frac{|m_\nu|\, y_N\, M_N^2}{2\sqrt{2}\,\pi^2\, v_h\,(m_{H_p}^2 - m_H^2)} \simeq -\frac{4.66\times 10^{-11}}{1-(m_{H_p}/125\ \text{GeV})^2}\, y_N \left(\frac{|m_\nu|}{0.05\ \text{eV}}\right)\left(\frac{M_N}{10\ \text{TeV}}\right)^2$$

M. DI Mauro. arXiv:2511.19622

Mediator decays before BBN



$$\tau_{\rm ferm}(H_p) \simeq \left(\frac{1.20\ {\rm s}}{y_N^2}\right)\left(\frac{20\ {\rm GeV}}{m_{H_p}}\right)\left(\frac{16\ {\rm TeV}}{M_N}\right)^4\left(\frac{0.05\ {\rm eV}}{m_\nu}\right)^2\left[1-\left(\frac{m_{H_p}}{125\ {\rm GeV}}\right)^2\right]^2$$

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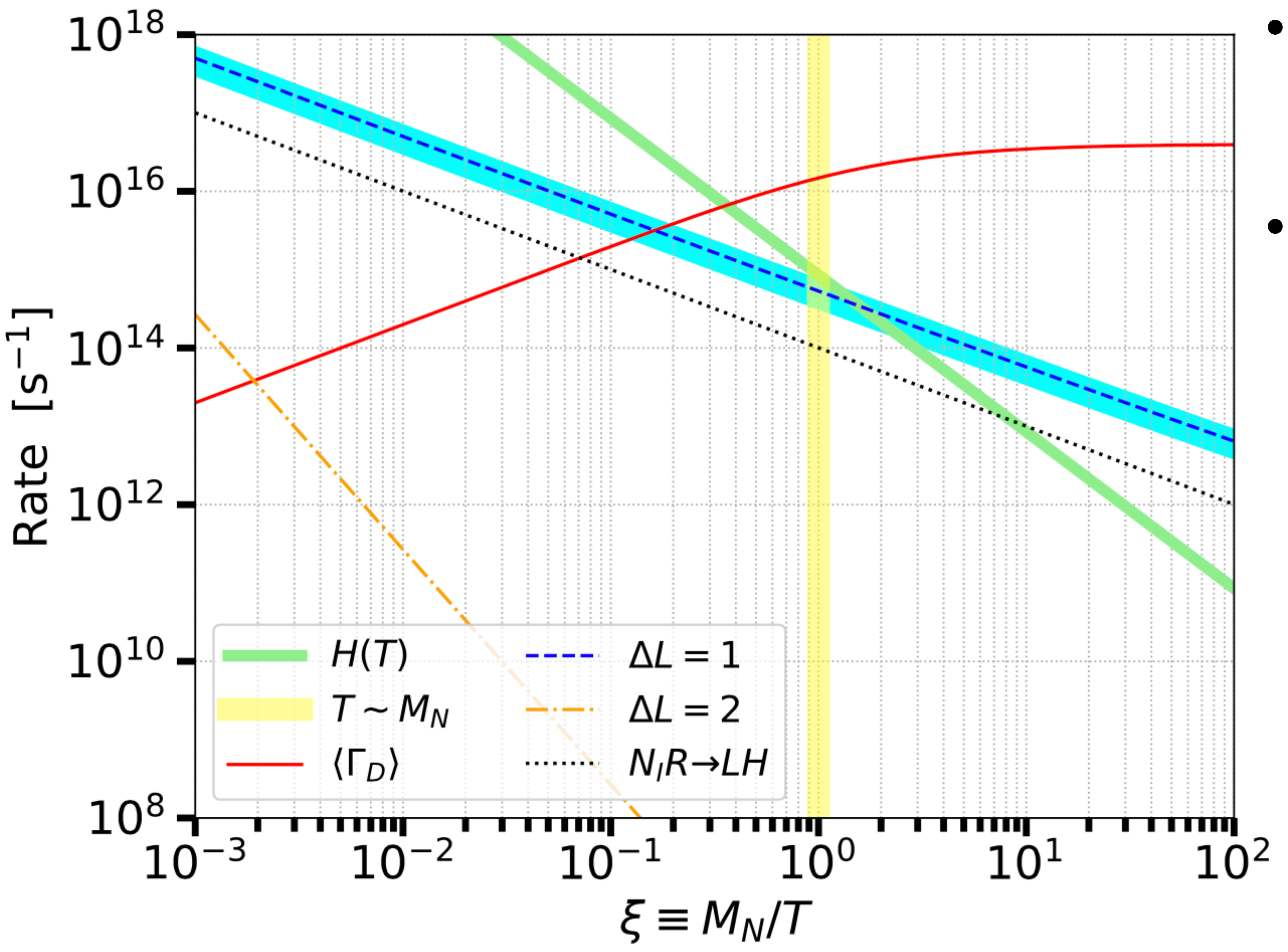
Dark Matter freeze out



- As the temperature drops below MN,I its abundance becomes Boltzmann suppressed and portal–mediated energy exchange shuts off.
- From that moment on, the comoving entropies of the two sectors are separately conserved

$$\zeta(T) = \left[\frac{g_{*S}^{\rm vis}(T)}{g_{*S}^{\rm vis}(T_{\rm dec})}\right]^{1/3} \left[\frac{g_{*S}^{\rm dark}(T'_{\rm dec})}{g_{*S}^{\rm dark}(T')}\right]^{1/3} \equiv T'/T$$

$$\zeta \simeq 1\text{–}2$$

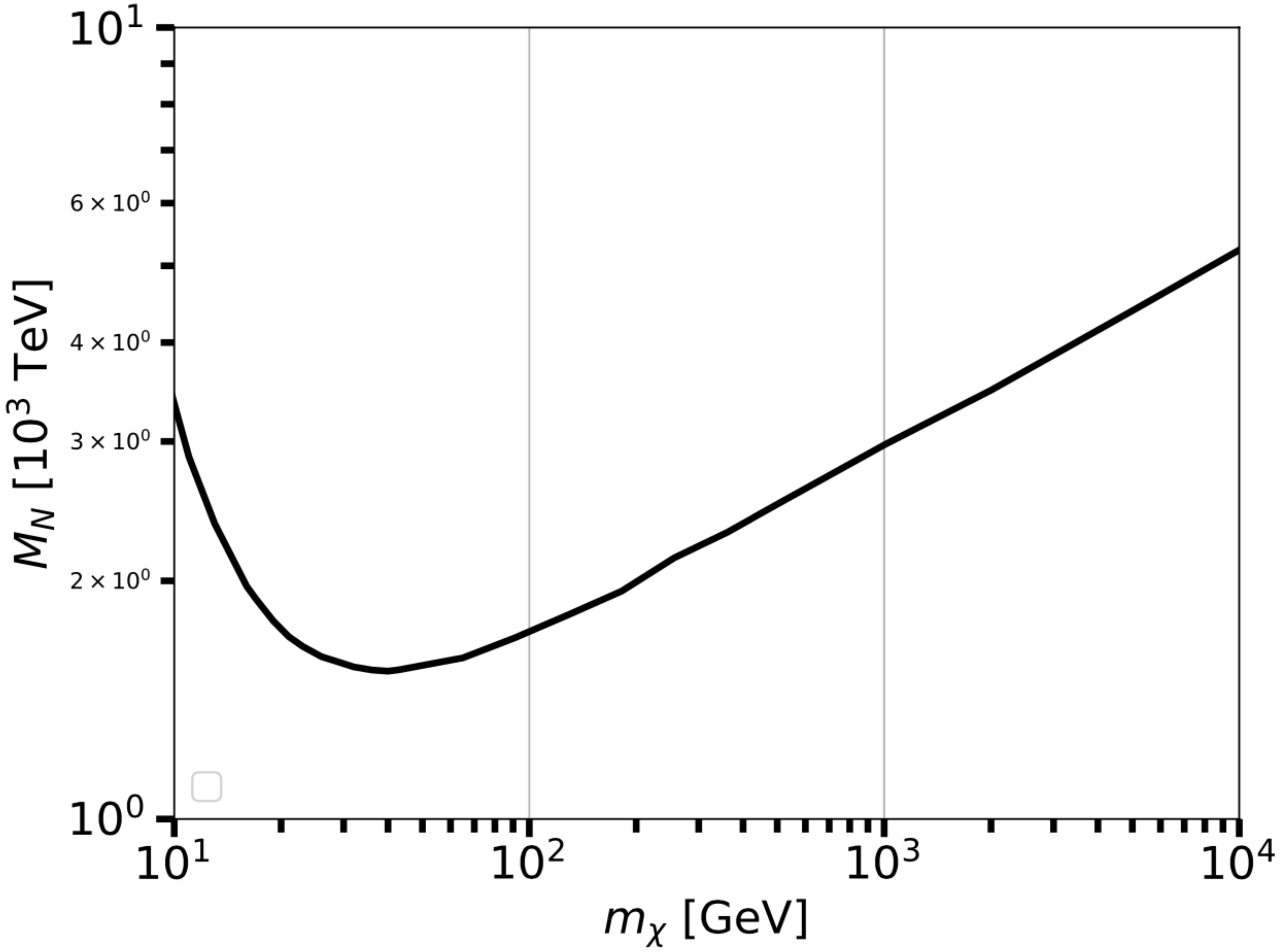
$$\Omega_\chi h^2 \simeq \left(1.05\times 10^9\ {\rm GeV}^{-1}\right) \frac{\sqrt{g_*^H(T_f, T'_f)}}{g_{*S}^{\rm vis}(T_f)} \frac{\zeta\, x_f'^2}{M_{\rm Pl}} \frac{64\pi m_\chi^2}{3\, y_p^4}$$

$$\mathcal{L}_{H,H_p-\chi\bar\chi} = -\frac{y_p}{\sqrt{2}}\left(\sin\alpha\, H + \cos\alpha\, H_p\right)\bar\chi\chi$$

$$y_p \simeq 0.43 \left(\frac{x'_f}{25}\right)^{1/2} \left(\frac{m_\chi}{100\ {\rm GeV}}\right)^{1/2} \left(\frac{g_{*S}^{\rm vis}}{86.25}\right)^{-1/4} \left(\frac{g_*^H}{100}\right)^{1/8},$$

M. DI Mauro. arXiv:2511.19622

Bounds from direct detection



$$M_N \lesssim 1.7{\times}10^3 \text{ TeV}\left(\frac{1}{y_p y_N}\right)^{1/2}\left(\frac{0.05 \text{ eV}}{m_\nu}\right)^{1/2}\left(\frac{m_{H_p}}{10 \text{ GeV}}\right)$$

M. DI Mauro. arXiv:2511.19622

Conclusions

- Direct, indirect, and collider searches are pushing standard WIMP models into increasingly small regions of parameter space.
  - In particular, scenarios where the same couplings govern both the relic density and direct-detection signals are now strongly constrained.
- **Secluded DM models** offer an elegant alternative, since the thermal history is decoupled from laboratory signatures.
- In these models, the relic abundance is set by interactions with a mediator that was in thermal equilibrium with the SM in the early Universe, while the portal to the visible sector remains very small.
- Such tiny couplings — or even a vanishing tree-level portal — can naturally arise in sequestered setups or through an appropriate symmetry.
- ***The portal then appears only at loop level, and heavy neutrinos above the TeV scale can simultaneously explain the smallness of neutrino masses and the invisibility of DM in laboratory searches.***

Backup slides

Dark matter with Beyond the standard Models

Dark Matter simplified Models

$$\mathcal{L}^{\chi}_{\rm int} = \xi\, \mu_\chi \lambda_\chi\, \chi^2 S + \xi\, \lambda_\chi^2\, \chi^2 S^2 + \sum_f \frac{g_f}{\sqrt{2}} \frac{m_f}{v_h}\, \bar{f} f\, S,$$

$$\mathcal{L}^{\psi}_{\rm int} = \xi\, \lambda_\psi\, \bar{\psi}\psi\, S + \sum_f \frac{g_f}{\sqrt{2}} \frac{m_f}{v_h}\, \bar{f} f\, S,$$

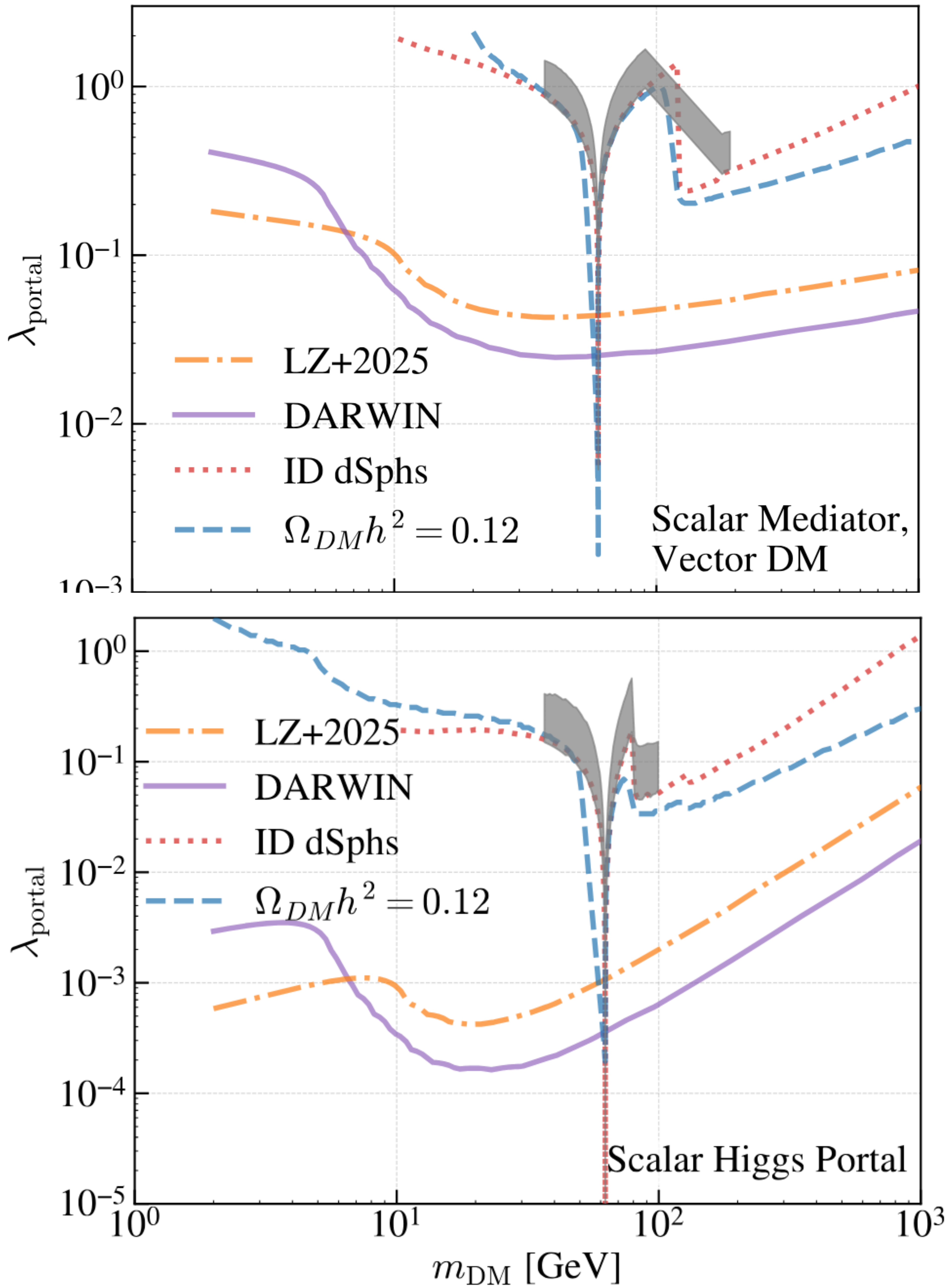
$$\mathcal{L}^{V}_{\rm int} = \mu_V \lambda_V\, V_\mu V^\mu S + \tfrac{1}{2}\lambda_V^2 V_\mu V^\mu S^2 + \sum_f \frac{g_f}{\sqrt{2}} \frac{m_f}{v_h}\, \bar{f} f\, S,$$

Higgs Portal models

$$\mathcal{L}_\chi = \tfrac{1}{2}(\partial_\mu \chi)^2 - \tfrac{1}{2}\mu_\chi^2 \chi^2 - \tfrac{1}{4!}\lambda_\chi \chi^4 - \tfrac{1}{2}\lambda_{h\chi}\chi^2 H^\dagger H,$$

$$\mathcal{L}_V = -\tfrac{1}{4}W_{\mu\nu}W^{\mu\nu} + \tfrac{1}{2}\mu_V^2 V_\mu V^\mu - \tfrac{1}{4!}\lambda_{hV}(V_\mu V^\mu)^2 + \tfrac{1}{2}\lambda_{hV} V_\mu V^\mu H^\dagger H,$$

$$\mathcal{L}_\psi = \bar{\psi}(i\not{\partial} - \mu_\psi)\psi - \frac{\lambda_{h\psi}}{\Lambda_\psi}\, \bar{\psi}\psi\, H^\dagger H,$$



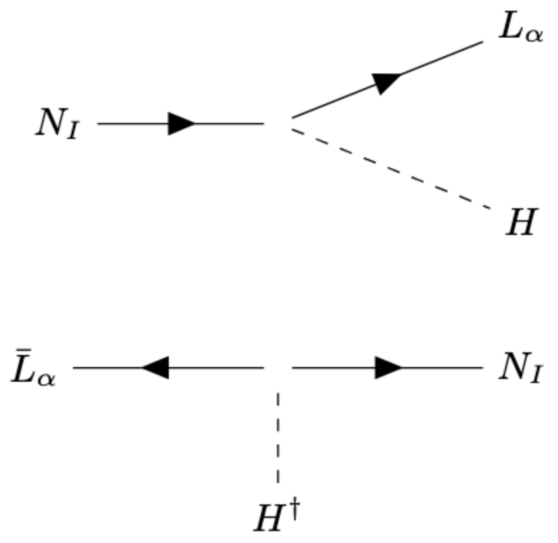
Kong, Di Mauro arXiv:2511.21808

Neutrino portal processes

(i) *Decays and inverse decays* (with violation of lepton number $\Delta L = 1$):

$$N_I \leftrightarrow L_\alpha H, \qquad N_I \leftrightarrow \bar{L}_\alpha H^\dagger. \tag{B3}$$

These control the production and destruction of heavy neutrinos N_I and are the leading neutrino–portal processes in the relativistic and mildly non–relativistic regimes.

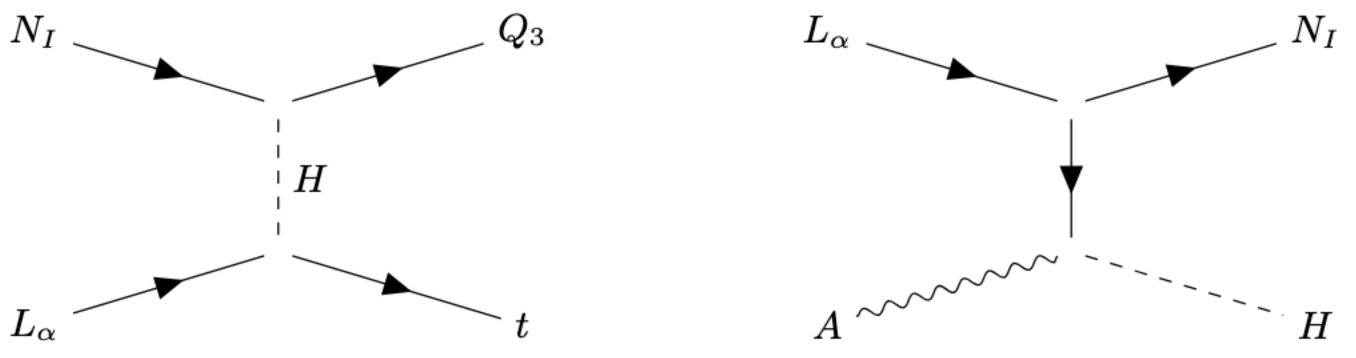


(*i*) Decays and inverse decays $N_I \leftrightarrow L_\alpha H,\ \bar{L}_\alpha H^\dagger$

(ii) $\Delta L = 1$ *Yukawa–mediated scatterings* with one Y_ν and one SM coupling, e.g.

$$N_I L_\alpha \leftrightarrow Q_3 t, \qquad L_\alpha A \leftrightarrow N_I H, \tag{B4}$$

and their crossed channels, mediated by Higgs or lepton exchange, where A is a SM electroweak gauge boson. These are $\mathcal{O}(Y_\nu^2 y_t^2)$ and $\mathcal{O}(Y_\nu^2 g^2)$ processes and efficiently contribute to keeping N_I in kinetic and chemical contact with the SM plasma for $T \gtrsim M_I$.

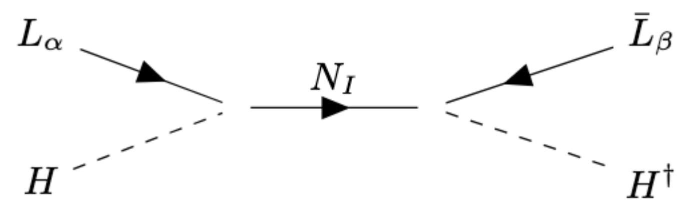


(*ii*) $\Delta L = 1$ scatterings: $N_I L_\alpha \leftrightarrow Q_3 t$ and $L_\alpha A \leftrightarrow N_I H$ (t-channel)

(iii) $\Delta L = 2$ *scatterings* induced by virtual N_I, such as

$$L_\alpha H \leftrightarrow \bar{L}_\beta H^\dagger, \qquad L_\alpha L_\beta \leftrightarrow HH, \tag{B5}$$

which encode lepton–number violation and are relevant for washout and for the low-energy Weinberg operator once N_I are integrated out. These processes are relevant only at very high temperatures ($T \gtrsim 10^{12}$ GeV) and are negligible near (and below) $T \sim M_I$. We will not consider them further.

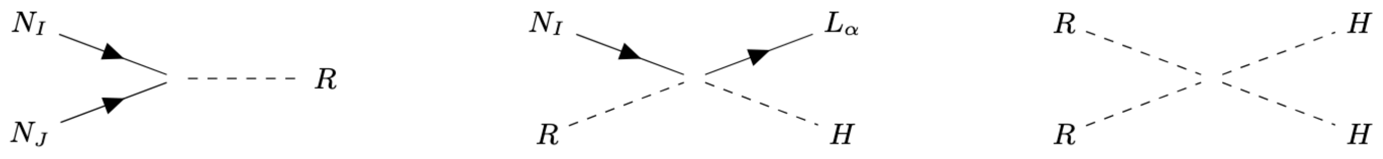


(*iii*) $\Delta L = 2$ scatterings: $LH \leftrightarrow \bar{L}H^\dagger$, $LL \leftrightarrow HH$ via s-channel N_I

(iv) *Portal processes involving the dark scalar and DM*, mediated by N_I and R, e.g.

$$N_I N_J \leftrightarrow R, \qquad N_I R \leftrightarrow L_\alpha H, \qquad RR \leftrightarrow HH, \tag{B6}$$

and, via the R–χ coupling, reactions that transfer energy and number between χ and the SM through the N_I–R chain. Once N_I are in equilibrium with the SM, these interactions help ensure that the dark sector is also thermally linked.



(*iv*) Portal processes with R (and χ via R): $N_I N_J \leftrightarrow R$, $N_I R \leftrightarrow L_\alpha H$, $RR \leftrightarrow HH$

Decoupling of NI

As the temperature drops below the mass of the lightest heavy neutrino $M_{N,I}$, the heavy–neutrino abundance becomes Boltzmann suppressed and portal–mediated energy exchange shuts off. From that moment on, the comoving entropies of the two sectors are separately conserved,

$$s_{\rm vis}(T)\,a^3 = \text{const}, \qquad s_{\rm dark}\big(T', \{\mu_i(T')\}\big)\,a^3 = \text{const}, \tag{C1}$$

where μ_i is the chemical potential for the particle species i. The temperature ratio at later times is fixed by separate entropy conservation after decoupling. Assuming full chemical equilibrium in the dark sector at the decoupling temperature $T'_{\rm dec}$ ($\mu_i(T'_{\rm dec}) = 0$), the general expression is (see, e.g., [56])

$$\zeta(T) \equiv \frac{T'}{T} = \left[\frac{s_{\rm vis}(T)}{s_{\rm vis}(T_{\rm dec})}\right]^{1/3} \left[\frac{s_{\rm dark}(T'_{\rm dec}, \{\mu_i{=}0\})}{s_{\rm dark}(T' = \zeta T, \{\mu_i(T')\})}\right]^{1/3}. \tag{C2}$$

If the following two physical conditions apply: (i) the relevant species in each sector behave as a radiation bath so that $s = \frac{2\pi^2}{45} g_{*S} T^3$, and (ii) *number–changing* reactions are fast enough to enforce $\mu_i = 0$, Eq. (C2) reduces to the familiar result (see e.g. [56])

$$\zeta(T) = \left[\frac{g_{*S}^{\rm vis}(T)}{g_{*S}^{\rm vis}(T_{\rm dec})}\right]^{1/3} \left[\frac{g_{*S}^{\rm dark}(T'_{\rm dec})}{g_{*S}^{\rm dark}(T')}\right]^{1/3}. \tag{C3}$$

In typical benchmarks with $m_\chi \sim \mathcal{O}(\text{GeV})$ and $M_I \sim \mathcal{O}(\text{TeV})$, the ensuing ζ lies in the ballpark $\zeta \simeq 1$–2 (the precise value follows from the g_{*S} evolution in the two sectors).

Dark sector freeze-out

After visible–dark decoupling, the portal remains too feeble to re–equilibrate the sectors. The dark bath stays internally thermalized at temperature $T' = \zeta T$ through fast secluded reactions. For the Dirac dark matter particle χ (mass m_χ) annihilating into a lighter real scalar H_p ($m_{H_p} < m_\chi$) via the Yukawa interaction $y_p \bar{\chi}\chi H_p$, the dominant number–changing channel is

$$\chi\bar{\chi} \to H_p H_p. \tag{C4}$$

We work with the visible–entropy yield $Y_\chi \equiv n_\chi / s_{\rm vis}$, using $s_{\rm vis}(T) = \frac{2\pi^2}{45} g_{*S}^{\rm vis}(T)\, T^3$, and define

$$x \equiv \frac{m_\chi}{T}, \qquad x' \equiv \frac{m_\chi}{T'} = \frac{x}{\zeta}. \tag{C5}$$

During radiation domination, the Hubble rate receives contributions from the radiation in both sectors,

$$H(T) = \sqrt{\frac{8\pi G}{3}\left(\rho_{\rm vis}(T) + \rho_{\rm dark}(T')\right)} = \frac{\pi}{\sqrt{90}}\frac{T^2}{M_{\rm Pl}}\left[g_*^{\rm vis}(T) + g_*^{\rm dark}(T')\,\zeta^4\right]^{1/2} \equiv \frac{1.66\sqrt{g_*^H(T,T')}}{M_{\rm Pl}}\,T^2, \quad ($$

The Boltzmann equation for Y_χ reads

$$\frac{dY_\chi}{dx} = -\frac{s_{\rm vis}(T)}{x\,H(T)}\langle\sigma v\rangle_{\chi\bar{\chi}\to H_p H_p}(x')\left[Y_\chi^2 - (Y_\chi^{\rm eq}(x'))^2\right], \tag{C8}$$

with $Y_\chi^{\rm eq}(x') = n_\chi^{\rm eq}(T')/s_{\rm vis}(T)$ and where the annihilation rate must be evaluated at the *dark* temperature T'.

$$y_p \simeq 0.43\left(\frac{x'_f}{25}\right)^{1/2}\left(\frac{m_\chi}{100\ {\rm GeV}}\right)^{1/2}\left(\frac{g_{*S}^{\rm vis}}{86.25}\right)^{-1/4}\left(\frac{g_*^H}{100}\right)^{1/8} \qquad \langle\sigma v\rangle_{\rm req} \approx \frac{2\times10^{-26}\ {\rm cm^3\,s^{-1}}}{\zeta}\left[1 + \mathcal{O}\!\left(\frac{g_*^{\rm dark}}{g_*^{\rm vis}}\zeta^4\right)\right]$$

Zero tree level portal

- The scalar portal (and similarly the abelian kinetic mixing) is gauge invariant and is therefore generically present in a 4D renormalizable EFT unless it is forbidden by additional UV structure.
- In our framework we impose the boundary condition.

$$\kappa(\Phi^\dagger\Phi)R^2 \qquad\qquad \kappa(\Lambda)=0$$

$$\beta_\kappa(\mu) \equiv \frac{d\kappa}{d\ln\mu} = \kappa(\mu)\,\mathcal{F}(\lambda_H, \lambda_R, y_p, g_i, \ldots),$$

- **Sequestered sectors**: A simple way to forbid all renormalizable tree-level portals is to assume that, at the UV scale Λ, the visible and dark sectors are geometrically or dynamically separated so that the renormalizable Lagrangian factorizes.
- **Discrete symmetry:** A complementary (and often orthogonal) protection is a discrete symmetry in the dark sector that forbids the portal at tree level.

M. DI Mauro. arXiv:2511.19622

Secluded vs WIMP

a. WIMP-like regime ($m_\chi < m_{H_p}$). The WIMP regime applies for $m_\chi < m_{H_p}$ with $y_p \sim \mathcal{O}(0.1$–$1)$ and a sizable scalar mixing $\sin\alpha \sim \mathcal{O}(0.1)$. In this case, annihilation into SM states is efficient because the dominant channels are $\chi\bar{\chi} \to f\bar{f}$ via s-channel exchange of H and H_p,

$$\langle\sigma v\rangle_{f\bar{f}} \propto y_p^2 \big(\sin\alpha\cos\alpha\big)^2, \tag{27}$$

so the same portal that controls annihilation also controls interactions with nuclei. For weak-scale masses, values of $y_p \sim 0.1$–1 can readily reproduce the observed relic density through thermal freeze-out. However, because the SI scattering amplitude is also proportional to the Higgs–singlet mixing, this region is strongly constrained by direct detection. A well-known exception occurs near the resonant regime, $m_\chi \simeq m_{H_p}/2$ (or $m_\chi \simeq m_H/2$), where annihilation is enhanced and the required couplings can be reduced [26], at the price of some mass tuning.

b. Secluded regime ($m_\chi > m_{H_p}$). In the secluded case, valid for $m_\chi > m_{H_p}$ with moderately small $y_p \sim \mathcal{O}(0.1)$ and $\sin\alpha \ll 1$, the relic abundance is set primarily by annihilations into dark-sector states, most importantly $\chi\bar{\chi} \to H_pH_p$ through t/u-channel χ exchange. For a Dirac fermion with a scalar Yukawa interaction, this channel is p-wave suppressed, so parametrically

$$\sigma v_{\rm rel}(\chi\bar{\chi} \to H_pH_p) = \frac{y_p^4}{64\pi m_\chi^2}\,\frac{\sqrt{1-r}}{\left(1-\frac{r}{2}\right)^4}\left(1 - r + \frac{r^2}{8}\right) v_{\rm rel}^2 \tag{28}$$

where $r = m_{H_p}^2/m_\chi^2$ and $v_{\rm rel}$ is the DM relative velocity.

This has two important implications. First, the freeze-out process can still be efficient because $v_{\rm rel}^2$ at freeze-out is $\mathcal{O}(0.1)$, whereas today in the Galactic halo $v_{\rm rel}^2 \sim 10^{-6}$, so late-time annihilation signals are further suppressed and indirect-detection constraints are typically very weak. Second, since the dominant annihilation does not rely on the portal to SM fields, the relic density can be achieved even when the mixing angle is extremely small.

SI scattering on nucleons proceeds through t-channel exchange of H or H_p and scales as (see Sec. VII)

$$\sigma_{\chi N}^{\rm SI} \propto y_p^2 \cos^2\alpha\,\sin^2\alpha \;\propto\; y_p^2\big(\tan 2\alpha\big)^2, \tag{29}$$

so in the secluded regime the nuclear cross section is naturally suppressed by $\sin^2\alpha$. Likewise, annihilation into SM fermions (Eq. 27) becomes negligible for $\sin\alpha \ll 1$, strongly weakening indirect-detection signatures associated with SM final states. For $y_p \sim \mathcal{O}(0.1$–$1)$, the process $\chi\bar{\chi} \to H_pH_p$ can nonetheless yield $\Omega_{\rm DM}h^2 \simeq 0.12$ (see Sec. VI).