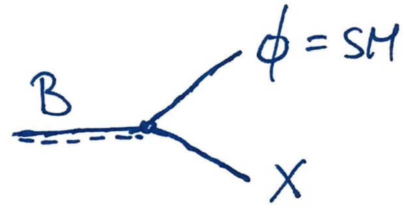
FREEZE - IN [0911.1120]

•) The freeze-in mechanism considers a very weakly coupled particle X, which is coupled to the heat bath B via λ. Suppose we have a "heavy" particle in the bath B, which then decays to X & the SM ϕ.

TOY MODEL e.g.: DECAY ONLY $\lambda \bar{X} B \phi$



$$\Gamma_B = \frac{\lambda^2}{8\pi} m_B$$

•) Naive estimate : $\frac{dn}{n} \sim \Gamma dt$, $H \propto T^2 \sim \frac{1}{t}$ for rad. dom.

$$Y_X = \frac{n_X}{s} \simeq \gamma_B \Gamma_B^0 \cdot t \simeq \frac{m_B}{T} \Gamma_B^0 \frac{1}{H} \simeq \frac{m_B \Gamma_B^0}{T^3} \simeq \frac{(m_B \lambda)^2}{T^3}$$

•) The yield of X depends on m_B & λ and shuts off at $T \to \infty$, dominated by low T.

•) A bit more rigorously:

$$a^{-3}\frac{d(n_X a^3)}{dt} = \int_{\Pi_{XB\phi}} \overline{|\mathcal{M}|^2} f_B$$

we neglected the "inverse decay" due to the lack of X in the plasma and $(1\pm f_X)\to 1$, $(1\pm f_{SM})\to 1$

$$\overset{\substack{0 \\ \shortparallel}}{}$$

$$\simeq 2 g_B \int_{\Pi_B} \Gamma_B m_B f_B$$

$$= \cancel{2} g_B \int \frac{d^3 p_B}{(2\pi)^3 \cancel{2} E_B} m_B \Gamma_B f_B$$

(dropping the index, all B)

$p\,dp = E\,dE$

$$= g \int_0^\infty \frac{4\pi}{8\pi^3} p^2 dp \frac{e^{-E/T}}{E} m \Gamma$$

$\frac{E}{m} = y, \; x = \frac{m}{T}$

$$= \frac{g m \Gamma}{2\pi^2} \int_m^\infty \sqrt{E^2 - m^2}\, e^{-E/T} dE$$

$[dE = T dx$

$$= \frac{g m \Gamma}{2\pi^2} \int_{\frac{m}{T}}^\infty \sqrt{x^2 - \left(\frac{m}{T}\right)^2} e^{-x} dx \, T^2]$$

$$* \; K_n(x) = \frac{\sqrt{\pi}}{(n-\frac{1}{2})!}\left(\frac{x}{2}\right)^n \int_1^\infty e^{-xy}(y^2-1)^{n-1/2} dy$$

$$= \frac{g m \Gamma}{2\pi^2} \underbrace{\int_1^\infty m \sqrt{y^2-1}\, e^{-xy} dy \, m}$$

$$= \frac{g m^3 \Gamma}{2\pi^2} \frac{K_1(x)}{x} = \frac{g m^2 \Gamma}{2\pi^2} T K_1\left(\frac{m}{T}\right)$$

·) We thus have:

$$a^{-3}\frac{d(n_x a^3)}{dt} = \frac{g m^2 \Gamma}{2\pi^2} T K_1\left(\frac{m}{T}\right)$$

Introducing the usual yield: $Y = \frac{n}{s}$, $s = \frac{2\pi^2}{45} g_{*s} T^3$

and $$a^{-3}\frac{d(Y_x s a^3)}{dt} = s\frac{dY_x}{dt} = s\frac{dY_x}{dT}\cdot\frac{dT}{dt}$$

radiation : $H = \frac{1}{2t}$, $\left(H^2 \propto a^{-4},\ \frac{a\,da}{} \propto dt,\ a \propto \sqrt{t},\ \frac{\dot a}{a} = \frac{1}{2t}\right)$

$$H \propto T^2 = \frac{1}{2t} \Rightarrow -\ln t = 2\ln T$$

$$-\frac{dt}{t} = 2\frac{dT}{T} \Rightarrow \frac{dT}{dt} = -HT$$

$$s\frac{dY_x}{dT}(-H\cancel{T}) = \frac{g m^2 \Gamma}{2\pi^2}\cancel{T} K_1\left(\frac{m}{T}\right)$$

$$Y_x = \frac{g m^2 \Gamma}{2\pi^2}\int_{T_{min}}^{T_{max}} \frac{K_1\left(\frac{m}{T}\right)}{sH} dT, \quad s = \frac{2\pi^2}{45} g_{*s} T^3,$$

$$H = 1.7\sqrt{g_*}\frac{T^2}{M_{Pe}}$$

$$= \frac{45}{1.7\ 4\pi^4}\frac{g\,\Gamma M_{Pe}}{m^2 g_{*s}\sqrt{g_*}}\underbrace{\int_{x_{min}}^{x_{max}} K_1(x)\, x^3 dx}$$

$$\int_0^\infty K_1 x^3 dx = \frac{3\pi}{2}$$

$$\simeq \frac{135}{1.7\ 8\pi^4}\frac{g_B}{g_{*s}\sqrt{g_*}}\left(\frac{M_{Pe}\Gamma_B}{m_B^2}\right)$$

·) The final abundance is

entropy density today

$$\Omega_X h^2 = \frac{\rho_X}{\rho_{cr}} = \frac{m_X n_X}{\rho_{cr}} = \frac{m_X Y_X s_0}{\rho_{cr}}$$

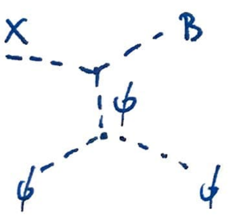
$$\rho_{cr} = \frac{3 H_0^2}{8\pi G}$$

$$= \frac{10^{27}}{g_{*s}\sqrt{g_*}} \left(\frac{m_X \Gamma_B}{m_B^2} \right)$$

·) To get the observed $\Omega_{DM} \sim 0.25$

$$\lambda \sim 10^{-13} \left(\frac{m_B}{m_X}\right)^{1/2} \left(\frac{g_*(m_B)}{100}\right)^{3/4} \left(\frac{g_B}{10^2}\right)^{-1/2}$$

·) In actual scenarios, one has to be careful to include other interactions and scatterings



$+ \ldots$ that may affect the picture.

Thermal freeze-in of sterile neutrinos

•) A well-motivated scenario for providing a DM candidate and neutrino masses, is the sterile neutrino ν_s.

ν_s is a neutral fermion, which is not coupled to the SM by gauge interactions, only via Yukawa.

$$\mathcal{L} = \mathcal{L}_{SM} + y_D \bar{L} \tilde{\Phi} \nu_s + M_s \nu_s^T C \nu_s \qquad L = \begin{pmatrix} \nu_L \\ \ell_L \end{pmatrix}$$

$$\phi = \begin{pmatrix} 0 \\ v \end{pmatrix}$$

in contrast to charged fermions, we can write down the Majorana mass term for ν_s. This is a Lorentz invariant term, which can be $M_s \gg M_W$.

$$\Rightarrow \mathcal{L}_{mass} \ni \left(\bar{\nu}_L \; \overline{\nu_S^c} \right) \begin{pmatrix} 0 & m_D \\ m_D & M_s \end{pmatrix} \begin{pmatrix} \nu_L \\ \nu_S^c \end{pmatrix}$$

From the 2x2 mass matrix, we get (6x6 for 3 generations)

$m_s \sim M_s$ ... the mass of ν_s

$m_\nu \sim - m_D^T M_s^{-1} m_D$ ... masses of light neutrinos.

•) The important point is that ν_s is very weakly coupled and is <u>not</u> thermalized in the early universe. After spontaneous breaking of $SU(2)_L$, when $\langle\phi\rangle = \begin{pmatrix} 0 \\ v \end{pmatrix}$ we have to rotate ν_L, ν_s into the mass basis:

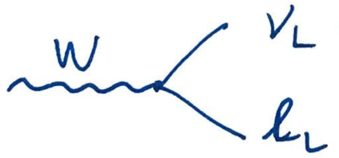
$$\nu_L^{(m)} \sim \nu_L + \theta \nu_s \quad ; \quad \theta \sim \frac{m_D}{M_s}$$
$$\nu_s^{(m)} \sim \nu_s - \theta \nu_L$$

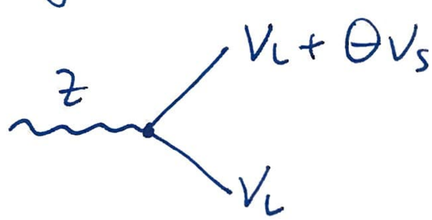
The mixing angle is very small $\theta \sim \sqrt{\underbrace{\frac{m_D^2}{M_s}}_{m_\nu} \frac{1}{M_s}} \sim \sqrt{\frac{m_\nu}{m_s}}$

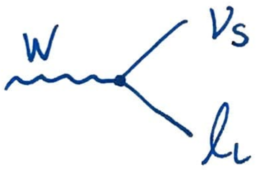
e.g. for $m_\nu \sim 10^{-4}\,\mathrm{keV}$ and $m_s \sim \mathrm{keV} \sim 10^{-6}\,\mathrm{GeV}$

$$\theta \sim \left(\frac{10^{-4}}{10^{3}}\right)^{1/2} \simeq 10^{-4} = \text{tiny}$$

•) Now we have small coupling to gauge bosons as well:





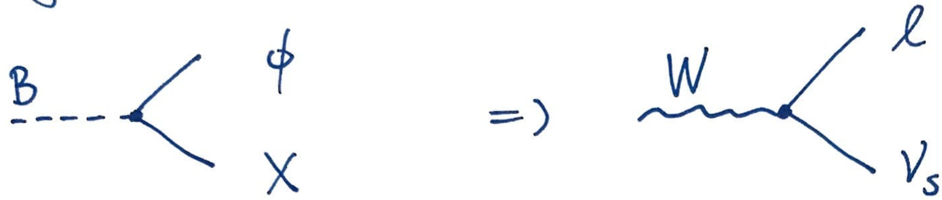
i.e.


$\propto \underline{\frac{g_L}{\sqrt{2}} \cdot \theta}\ \bar{l}_L \gamma^\mu \nu_s W_{L\mu}$

This leads to: $\langle \sigma v \rangle n_s \sim G_F^2 T^2 \theta^2 T^3 \sim G_F^2 T^5 \theta^2$

•) As in the freeze-out case above, the Y_{ν_s} starts to accumulate slowly as the temperature drops.

Using the same formalism as above



$$\Omega_{DM} h^2 = \frac{10^{27}}{g_{*s}\sqrt{g_*}\big|_{M_W}} \frac{m_s \Gamma_{W\to\nu_s e}}{M_W^2}$$

$$\Gamma_{W\to\nu_s e} = \frac{\alpha_W}{4}\theta^2 M_W \,, \quad g_*(M_W) \sim 100$$

$$\Omega_{\nu_s} h^2 = 0.1 \left(\frac{m_{\nu_s}}{\text{keV}}\right)\left(\frac{\theta}{10^{-7}}\right)^2$$

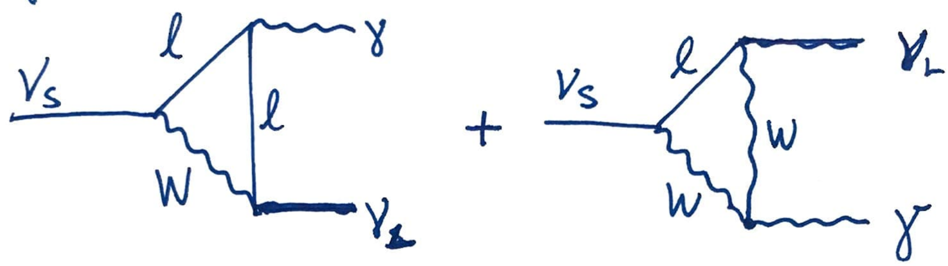
• There are important corrections in this scenario, such as ν-oscillations, T-dependent potentials

$$\theta_0 \to \theta(T) \simeq \frac{\theta_0}{1+\left(\frac{T}{T_0}\right)^6} \,, \quad T_0 \simeq 100\,\text{MeV}$$

Resonant production, refined estimate

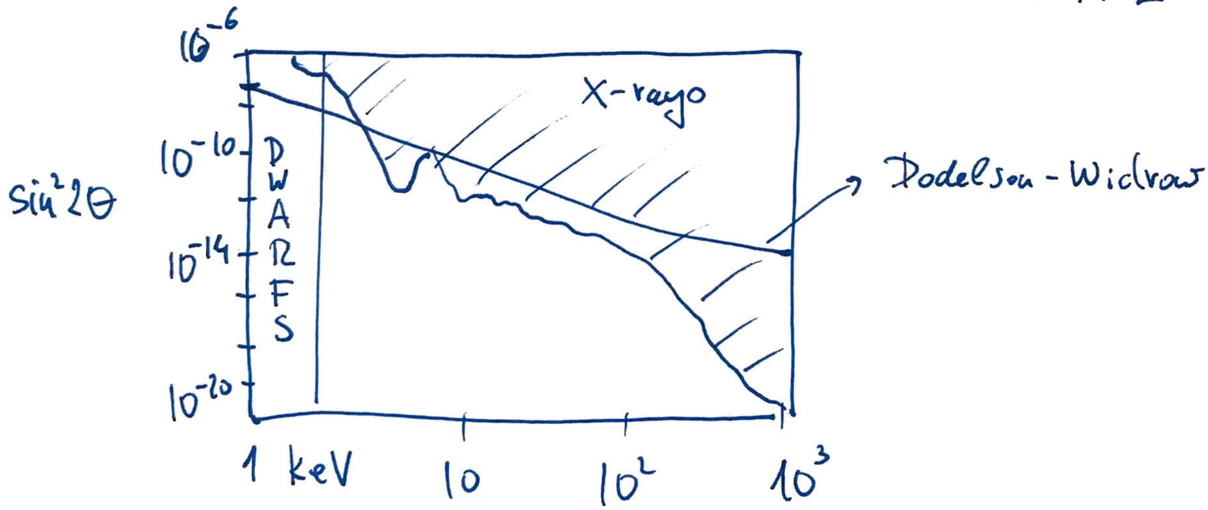
$$\Omega_{\nu_s} h^2 \simeq 0.1 \left(\frac{\theta^2}{3\cdot 10^{-9}}\right)\left(\frac{m_{\nu_s}}{3\,\text{keV}}\right)^{1.8}$$

•) This is referred to as the Dodelson-Widrow mechanism. It is strongly constrained by indirect searches. Because of the mixing, ν_s couples to the SM and is destabilized at one loop level: $\nu_s \rightarrow \gamma \nu_L$, $E_\gamma \sim m_{\nu_s}/2$



$$\tau = \frac{1}{\Gamma} = 10^{30}\,s \left(\frac{10^{-7}}{s^2_{2\theta}}\right)\left(\frac{keV}{m_{\nu_s}}\right)^5 \Rightarrow$$

X-ray with $E \sim keV$



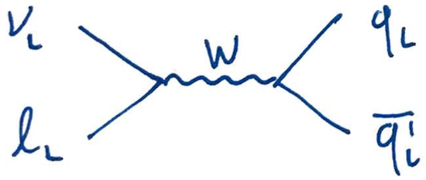
•) Presence of primordial lepton asymmetry or additional interactions opens up the allowed parameter space [1910.04901]

Thermal over-production and entropy dilution

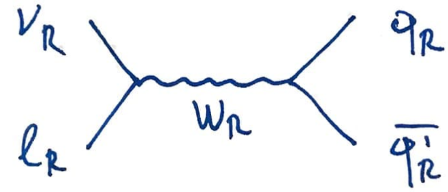
•) Let us re-consider the freeze-out of a massive but relativistic species. This is very similar to the SM neutrino decoupling but via BSM (weaker) interactions. A somewhat generic example is a W_R with $M_{W_R} \gg M_W \sim$ few TeV that couples to L_R, Q_R, This is motivated by parity restoration at high scales and necessarily contains the ν_R, which we take $m_{\nu_R} \sim$ keV.

•) Relativistic thermal freeze-out goes as in the SM.

WEAK



SUPER-WEAK



$$\Gamma \sim G_F^2 T^{\cancel{5}\,3} = \frac{\sqrt{g_*}\,\cancel{T^2}}{M_{Pl}}$$

$$\Rightarrow T_{\nu dec} \cong g_*^{1/6} \left(\frac{M_W^4}{M_{Pl}} \right)^{1/3}$$

$$G_F \propto \frac{1}{M_W^2} \rightarrow G_F' \propto \frac{1}{M_{W_R}^2}$$

$$G_F^2 \left(\frac{M_W}{M_{W_R}} \right)^4 T^3 = \frac{\sqrt{g_*}}{M_{Pl}}$$

$$T_f \sim 0.4\,\mathrm{GeV} \left(\frac{M_{W_R}}{5\,\mathrm{TeV}} \right)^{4/3}$$

•) With suppressed interactions, the freeze-out happens earlier $T_f \sim 0.4$ GeV, $g_*(T_f) \sim 70$.

Earlier decoupling means higher temperatures T_f GeV and ~~two~~ less Boltzmann suppression.

•) For $m_{\nu_R} \sim$ keV $\ll T_f \sim 0.4$ GeV, the number density is given by the relativistic formula for $m \ll T$

$$n_{\nu_R} \sim \overset{\overset{2}{\shortparallel}}{g} \frac{3}{2\pi^2} \zeta(3) T_f^3$$

$$Y_{\nu_R} = \frac{n_{\nu_R}}{s} = \frac{135\zeta(3)}{4\pi^4 g_*(T_f)} \quad , \; s = \frac{2\pi^2}{45} g_{*s} T^3$$

•) As for the usual SM neutrinos, the ν_R simply cools down and the energy density is

$$\rho_{\nu_R} = m_{\nu_R} n_{\nu_R} = m_{\nu_R} Y_{\nu_R} s$$

where $s_{today} \simeq \frac{2\pi^2}{45} \underbrace{\left(2 + 2 \cdot 3 \cdot \frac{7}{8} \left(\frac{T_\nu}{T_\gamma}\right)^3\right)}_{2+\frac{21}{4}\frac{4}{11} = \frac{43}{11}} T_\gamma^3 = \frac{2\pi^2}{45}\frac{43}{11} T_\gamma^3$

$\simeq 3000\ \text{cm}^{-1}$ [PDG '21]

•) The final energy density parameter is then given by

$$\Omega_{\nu_R} \simeq 3.3 \left(\frac{m_{\nu_R}}{\mathrm{keV}}\right)\left(\frac{70}{g_*(T_f)}\right)$$

phase-space bound

OVER-ABONDANT by

$$\frac{\Omega_{\nu_R}}{\Omega_{DM}} \sim \frac{3}{0.25} \sim 12.$$

•) This is quite similar to the SM with

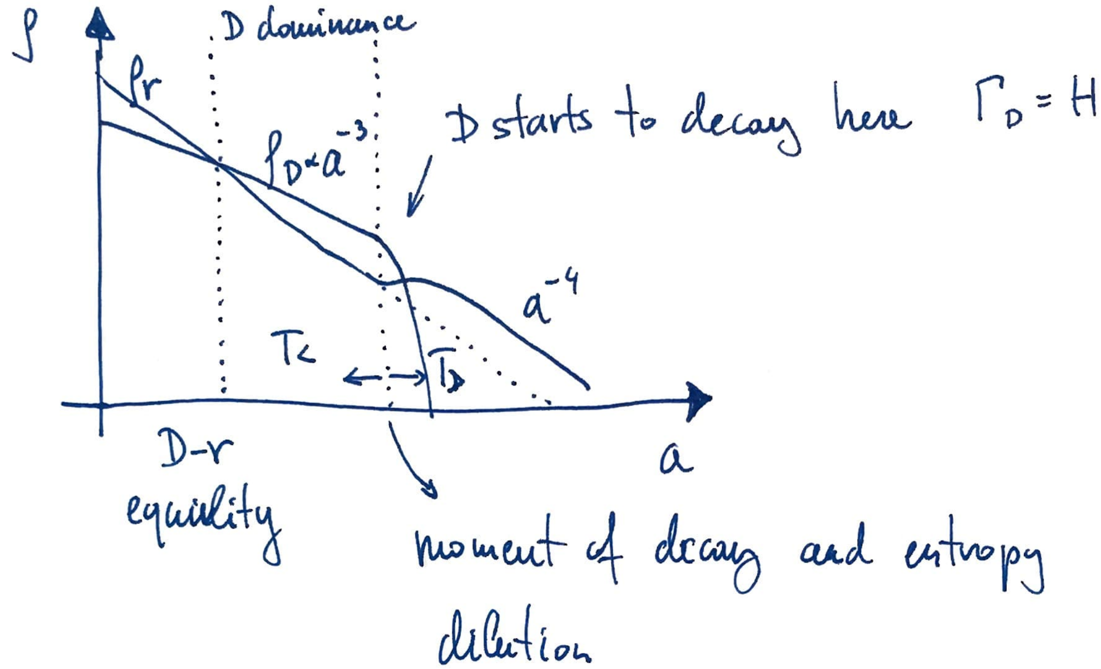
$$\Omega_{\nu_L} \simeq \frac{\sum m_{\nu_i}}{93\,\mathrm{eV}}$$

ENTROPY DILUTION BY LATE DECAYS [book by Kolb & Turner]

[Scherrer-Turner '85]

•) The issue of over-abundance (similar to the monopole problem) can be addressed by entropy production. This may happen when a massive species is long-lived enough to dominate the $\rho_m > \rho_r$ and then decays to finale states that thermalizes and go into ρ_r.

•) Let us refer to the massive species as the dilutor D with mass m_D & decay rate Γ_D.



•) Before the decay, D is matter and dominates

$$H^2 = \frac{8\pi G}{3}\rho_m = \frac{8\pi G}{3} Y_D m_D s = \frac{8\pi}{3 M_{Pl}^2} Y_D m_D \frac{2\pi^2}{45} g_{*s} T_<^3$$

$$= \Gamma_D^2 \Rightarrow T_< \sim \left(\frac{\Gamma_D^2 M_{Pl}^2}{Y_D m_D g_{*s}}\right)^{1/3}$$

•) After D decays into radiation, its $\rho_m \to \rho_r$ and the universe is again radiation-dominated.

$$\Gamma_D = H \simeq 1.7\sqrt{g_*}\frac{T_>^2}{M_{Pl}} \quad \text{or} \quad T_> = \sqrt{\frac{\Gamma_D M_{Pl}}{g_*^{1/2}}}$$

•) The ratio between these two is given by

$$\frac{T_>}{T_<} \simeq \frac{\sqrt{\Gamma M_{Pl}}}{(\Gamma M_{Pl})^{2/3}} \frac{(Y_D m_D g_*)^{1/3}}{g_*^{1/4}} = \left(\frac{g_*^{1/4} Y_D m_D}{\sqrt{\Gamma M_{Pl}}} \right)^{1/3}$$

Because the entropy goes as T^3, the dilution factor is given by

$$\frac{S_>}{S_<} = \frac{g_*^{1/4} Y_D m_D}{\sqrt{\Gamma_D M_{Pl}}}$$

•) Of course, this entropy is associated only to the SM particles (radiation), we assumed that D does <u>not</u> decay to ν_R. Thus the energy density $\rho_{\nu_R} = m_{\nu_R} n_{\nu_R}$ is diluted by precisely this factor. The takeaway message is that late decays heat up the radiation and effectively dilute the n of any particles that don't get produced by their decays.

•) SUMMARY : The amount of dilution depends on

* $Y_D m_D$: the more abundant Y_D is (e.g. if produced relativistically) the more it dilutes. Also, the higher its mass, the more ρ it carries and heats up radiation more.

* $\propto \Gamma_D^{-1/2}$ or $\tau_D^{1/2}$: the longer it lives, the more it dominates and the more effective the energy transfer = higher dilution.

Thermal overproduction $\qquad\qquad$ Entropy dilution

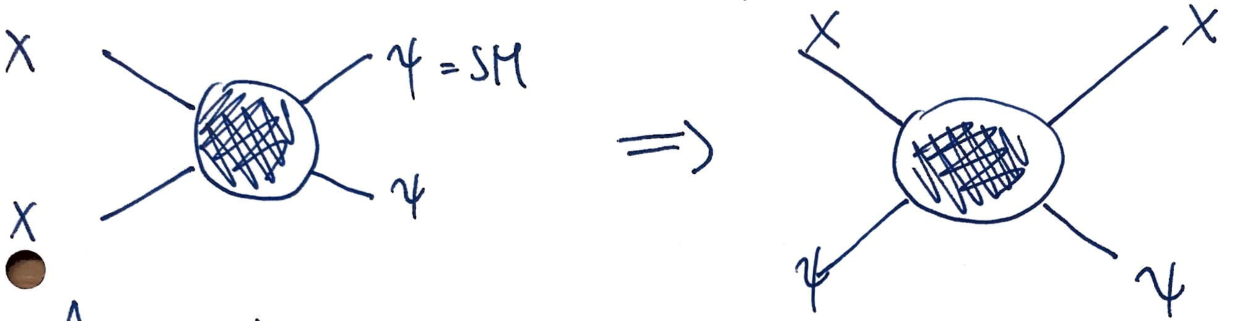
$$\Omega_{\nu_R} \gtrsim 3 \left(\frac{m_{\nu_R}}{\mathrm{keV}}\right) \Longrightarrow \Omega_{\nu_R} \sim 0.23 \left(\frac{m_{\nu_R}}{\mathrm{keV}}\right)\left(\frac{2\,\mathrm{GeV}}{m_D}\right)\left(\frac{1\mathrm{s}}{\tau_D}\right)^{1/2}$$

•) We will see, however, that the lifetime of D cannot be arbitrarily long $\tau_D \lesssim 1\mathrm{sec}$, otherwise it spoils the successful prediction of BBN in the SM.

DIRECT DETECTION OF DARK MATTER

[1904.07915]

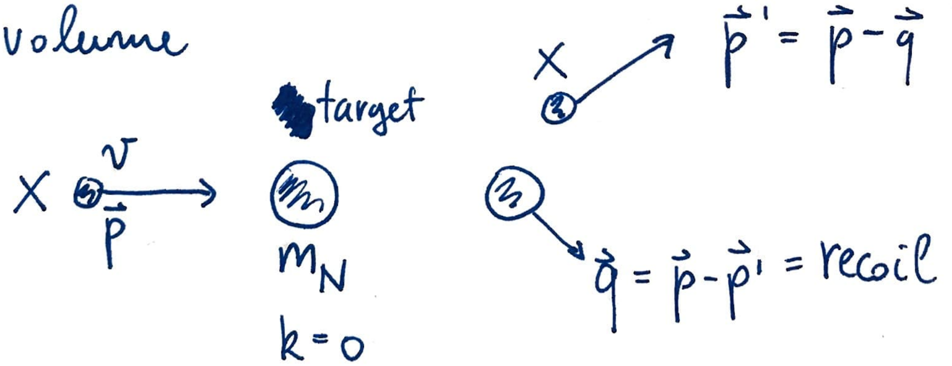
•) A typical WIMP scenario with $\langle \sigma v \rangle \sim 10^{-8}\,\text{GeV}^{-2}$ is an ideal candidate for direct detection.



ANNIHILATION
=
FREEZE-OUT

SCATTER OFF ψ
=
DIRECT DETECTION

•) The incoming DM with $v \sim 200\,\text{km/s}$ or $v/c \sim 10^{-3}$ can scatter off a nucleus sitting at rest in a target volume



•) The whole process is non-relativistic and thus

$$E_i = \frac{p^2}{2m_X}, \qquad E_f = \frac{(\vec{p}-\vec{q})^2}{2m_X} + \frac{q^2}{2m_N}$$

$$p = m_X v$$

•) Energy conservation will give us the max. recoil

$$E_f = E_i \Rightarrow \cancel{\frac{p^2 - 2pq c_\theta + q^2}{2m_X}} + \frac{q^2}{2m_N} = \cancel{\frac{p^2}{2m_X}}$$

$$\Rightarrow \frac{pq c_\theta}{m_X} = \frac{q^2 (m_X + m_N)}{2 M_X M_N} = \frac{q^2}{2\mu}$$

reduced DM-nucleon mass

Now since $|c_\theta| < 1$

$$q_{max} = \frac{2p\mu}{m_X} \simeq 2v\mu \simeq 10^{-3} \times (10-100)\,\text{GeV} \qquad m_N \sim A m_p$$

$\Downarrow$

maximum momentum transfer between X = DM & N.

•) From here, we get the largest recoil energy $E_{R_{max}}$

$$E_{R_{max}} = \frac{q^2_{MAX}}{2m_N} = \frac{2v^2\mu^2}{m_N} \sim 10^{-6} \frac{m_X^2 m_N}{(m_X + m_N)^2} \sim \underline{\underline{(10-100)\,\text{keV}}}$$

and: $v_{min} = \sqrt{\dfrac{E_R m_N}{2\mu^2}}$.

$\downarrow$

for a fixed measured E_R, this is the smallest v that DM can have to trigger a signal. Remember that v follows a Maxwell distribution and has suppressed high v tails $\propto e^{-(v/v_0)^2}$, $v_0 \sim 220 \frac{km}{s}$.

•) Furthermore, note that the momentum transfer

$$q \sim (10-100)\,\mathrm{MeV} \Rightarrow q^{-1} \sim \lambda \sim (1-10)\,\mathrm{fm},$$

which implies we are scattering DM on an <u>entire</u> nucleus, not individual p,n with $m_p \sim \mathrm{GeV}$.

•) The final number that we wish to measure is the rate of scattering / time / E_R : $\frac{dR}{dE_R}$. To calculate that, one has to go through a series of theories, none of which we can analyze in depth here (assuming QFT is not an option). So we will sketch the procedure in enough detail to appreciate the progress of direct detection exp.

DIRECT DETECTION THEORY LADDER

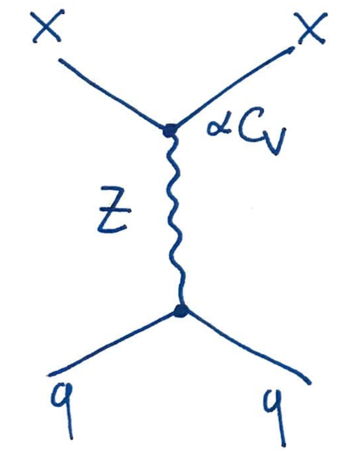
i) Start with $\mathcal{L}_{SM+DM}$; q's, ℓ's, W, Z, X (QCD)

ii, Match the theory to $\mathcal{L}_{p,n}$ = theory for non-relativistic neutrons, protons $q \sim 1$–$10\,\mathrm{MeV} \ll m_p$

iii) Calculate the amplitude for X–n couplings

iv) Compute the final X–nucleus N cross-section by averaging over the <u>local</u> DM velocity distribution.

strength of the DM-X, $g_2 \overset{?}{\sim} C_V$ or $C_V \ll 1$

i) 

: $\dfrac{C_V}{M_Z^2}\,\bar{X}\gamma^\mu X\,\bar{q}\gamma^\nu q$ ... vector coupling

spin-independent

$\dfrac{C_A}{M_Z^2}\,\bar{X}\gamma^\mu\gamma^5 X\,\bar{q}\gamma^\nu\gamma^5 q$

... axial coupling

spin-dependent

ii) This is now matched to n, p : $\frac{2}{m_N} \ll 1$

$$\langle n(k')|\,\bar{q}\gamma^\mu q\,|n(k)\rangle = \bar{u}_n(k')\left(F_1^n(q^2)\gamma^\mu + \frac{i}{2m_n}F_2^n(q^2)\sigma^{\mu\nu}q_\nu\right)u_n(k)$$

$u_n \to \sqrt{m_n}\,\xi_s$ $\quad$ $\bar{u}_n\gamma^0 u_n \to$

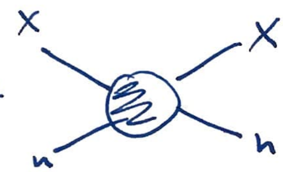
•) The $F_{1,2}(q^2)$ are the nucleon form-factors, just like the g-2 in QED, for example. These can be computed in the heavy baryon χPT.

iii) Taking the NR limit for both X ($v \sim 10^{-3}$) and for the n, p, the Z diagram gives

$$\mathcal{M} = \frac{C_V}{M_Z^2} m_X m_n \underbrace{\xi^+_{X'} \xi_X \; \xi^+_{n} \xi_n}_{\text{spinors for } X, n \;\Rightarrow\; \mathcal{O}(1)}$$

$$\overline{|\mathcal{M}|}^2 = \left(\frac{3C_V}{M_Z^2}\right)^2 \left(4 m_X m_n\right)^2$$

this is now integrated over the final 2-body phase-space of X and n



$$d\sigma_n = \frac{1}{4 m_X m_n v} \int_{\Pi_{X,n}} \overline{|\mathcal{M}|}^2 (2\pi)^4 \delta^{(4)}(\Sigma p) \quad , \quad \int_{\Pi_X} = \int \frac{d^3 p_X}{\underset{\underset{m_X}{\parallel}}{2E_X} (2\pi)^3}$$

$$\sigma_n = \frac{\mu^2}{\pi} \left(\frac{3C_V}{M_Z}\right)^2$$

iv) Finally, we go from nucleons n, p to nuclei

N = Xe, Ar, Si, Ge, ... with $A \sim 70-100$.

·) The differential σ is then given by
with $d|q|^2 = 2m_N dE_R$

known form-factor (Helm)

$$\frac{d\sigma_N}{dE_R} = \frac{\sigma_n m_N}{2(\mu v)^2} F^2(q)\,\Theta(v - v_{min}) \quad , \quad F^2 \sim A^2$$

$$\Downarrow$$

$$\frac{dR}{dE_R} = N_T n_X \int \underbrace{\frac{d\sigma_N}{dE_R} v f(v)}_{\langle \sigma v \rangle} d^3 v$$

$$\parallel$$

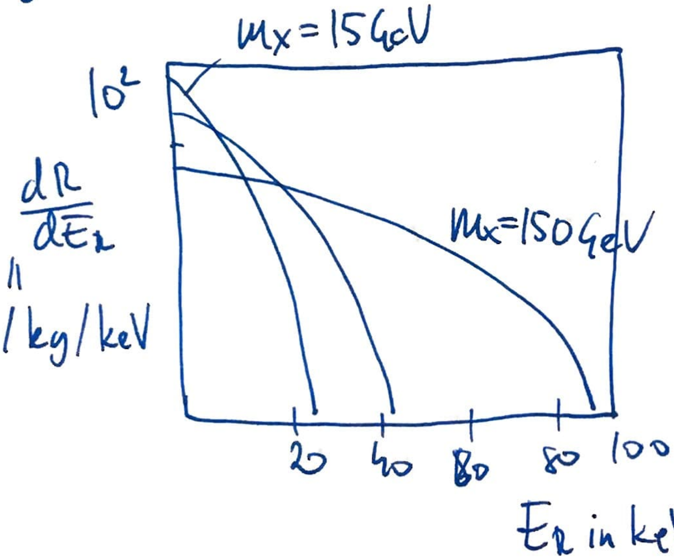
$$= N_T \frac{\rho_X}{m_X} \frac{\sigma_n m_N}{2\mu^2} F^2(q) \underbrace{\int \frac{f}{v} \Theta(v - v_{min}) d^3 v}_{g(v_{min})}$$

N_T: # of nuclei / mass
~ Avogadro × mass number

·) $g(v_{min}) \sim e^{-m_N E_R / 2\mu^2 v_0^2}$, $v_0 \sim 220$ km/s

$$F(q) = \frac{3}{q r_n} J_1(q r_n)\, e^{-(qs)^2/2} \quad , \quad r_n \sim \text{fm}\, A^{1/3} , \; s \sim \text{fm}.$$

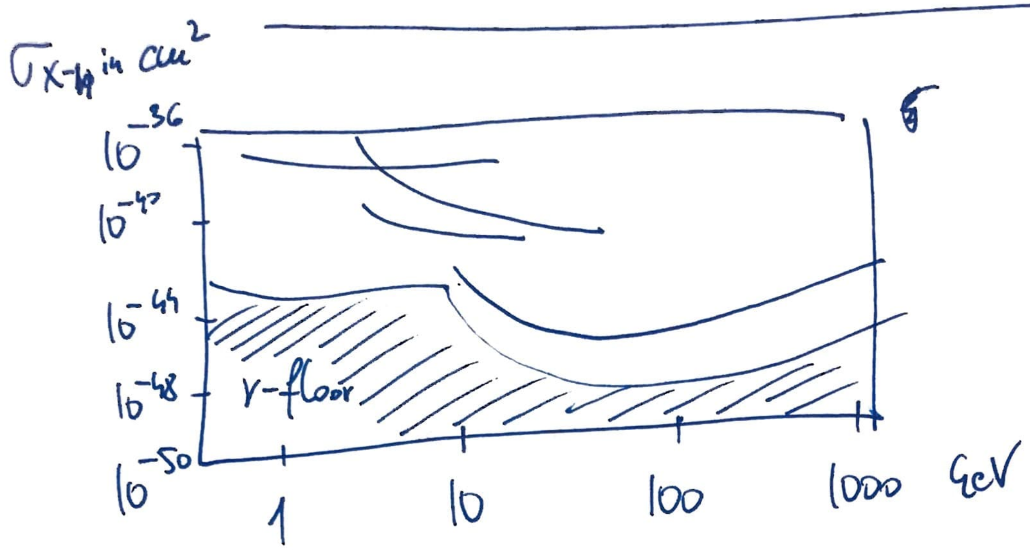
Helm f.f.



* the rate is exponentially suppressed for higher E_R: $g(v_{min}) \propto e^{-E_R}$

* better for low m_X, $v_{min} \propto \frac{1}{m_X}$

SUMMARY OF DIRECT DETECTION EXP.



*) At low m_X, the recoil is too small: $E_R \ll$ keV to be detectable by searches with $E_R^{thr} \sim$ few keV.

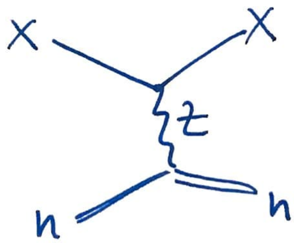
$\propto \mu^2 \propto m_X^2$.

*) At high m_X, the number density n_X goes down as m_X^{-1}, $\rho_X = m_X n_X$ = fixed by Ω_{DM}.

*) How strong are these limits?

taking $C_V \sim g_2 \sim 0.65$ or $O(1)$, we have

$$\sigma_n^V \simeq 10^{-37}\, cm^2$$



$m_X \sim 100$ GeV

This is way excluded $\Rightarrow$ needs $C_V \ll 1$, however this affects the freeze-out and WIMP "miracle starts to fade.

INDIRECT DETECTION [1710.05137]

•) The idea behind indirect detection is to look for visible products of DM <u>decay</u> or <u>annihilation</u>. The final states are SM particles such as γ, ν, $e, W, Z, p, n, \dots$).

The <u>upshot</u> for such searches is that the DM energy density is HUGE $\Omega_{DM} \cong 5\,\Omega_B$ so there is plenty of DM to produce a signal.

The <u>downsides</u> are:

i) DM might be very weakly coupled (perhaps even only gravitationally)

ii) There are many astrophysical backgrounds and uncertainties (like pulsars), which make it very hard to claim a signal from DM only.

Neutral final states propagate basically unimpeded (almost) like γ's and ν's. $\Rightarrow$ $(\theta, \varphi, \underset{\text{from the redshift}}{r})$

•) Charged final states ($e^\pm$, $p^\pm$, other cosmics) get deflected by ~~astro~~ intergalactic B-fields. This makes the analysis more involved (there are tools for this, like GALPROP) and harder to pinpoint the origin of the DM signal.

• <u>ESTIMATES</u> of number of events @ earth

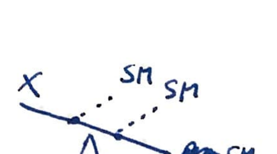
i) from DECAYS

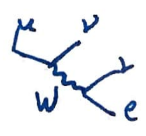
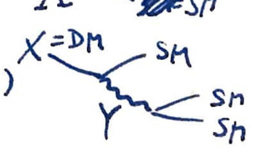
ii) from ANNIHILATIONS

i) As we saw in the case of νs, DM may decay

• A very generic estimate with $\mathcal{O}(1)$ couplings can be made on dimensional grounds $[\Gamma]=1$.

$$\Gamma = \begin{cases} \dfrac{M_X^3}{\Lambda^2} \sim 1\,s^{-1} & \text{(a)} \\[2ex] \dfrac{M_X^5}{\Lambda^4} \sim 10^{26}\,s^{-1} & \text{(b)} \end{cases} \qquad \text{where } \Lambda \sim M_{GUT} \sim 10^{16}\,\text{GeV}$$

(a) : This comes from a $d=5$ operator, e.g. 

(b) : From $d=6$ operator, like μ-decay  $\Rightarrow$ 

a) Is too fast $\Rightarrow$ DM unstable

b) Is long enough : $\tau_{DM} \sim 10^{26}\,s = 10^{19}\,yrs \gtrsim 10\,byrs$
$\sim t_0$

•) Clearly, it is very much feasible that DM is stable on cosmological scales and starts decaying at the present age.

• How many events can we expect to observe?

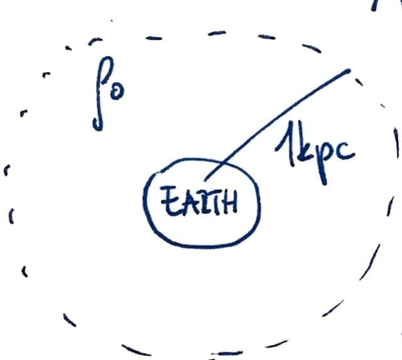
detector, surface $A \sim m^2$

dV, n_x ... local number density

r ... distance to the DM-rich object (dwSph, center of MW)

• $$\frac{dN}{dt} \simeq \underbrace{\frac{n_x\,dV}{\tau_x}}_{\propto \text{ source volume}} \underbrace{\frac{A}{4\pi r^2}}_{\text{geometric acceptance}}\,, \quad dV = r^2\,dr\,d\Omega$$

$$= A\,n_x \left(\frac{d\Omega}{4\pi}\right) \frac{dr}{\tau_x} \quad \text{(note } r^2 \text{ cancels away)}$$



$\rho_0 \sim 0.4\,GeV/cm^3$ $\Rightarrow$ $\dfrac{dN}{dt} = \dfrac{A n_x r}{\tau_x} = \dfrac{A \rho_0 r}{M_x \tau_x}$

$r \lesssim 1kpc$ for e^+e^- $E \sim M_W$ to propagate

$\Rightarrow$ a uniform local sphere with $r \sim 1\,\mathrm{kpc}$ around the earth gives us:

$$\frac{dN_{dec}}{dt} = 10^{-4}\,s^{-1} \left(\frac{A}{m^2}\right)\left(\frac{\rho_0}{0.4\,\frac{GeV}{cm^3}}\right)\left(\frac{r}{kpc}\right)\left(\frac{TeV}{m_X}\right)\left(\frac{10^{26}s}{\tau_X}\right)$$

This implies $\sim 10^3$ events per year.

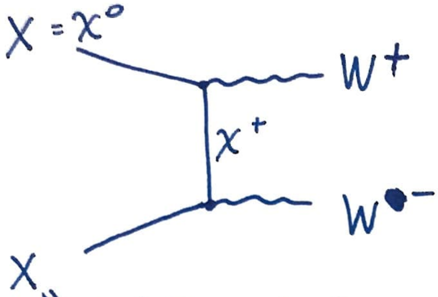
•) The case of sterile neutrino decay also gave a similar lifetime but $m_X \sim \mathrm{keV}$

$$\tau_X \sim 10^{30}\,s \left(\frac{10^{-7}}{\theta^2}\right)\left(\frac{keV}{m_X}\right)^5$$

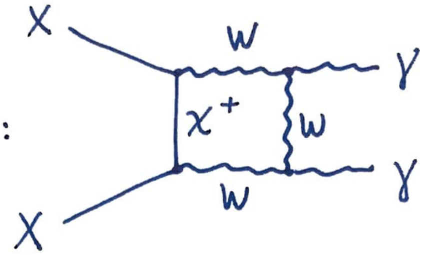
•) Even though the rate is higher, the total energy that goes into the signal is $\propto m\dot{N}$ and thus stays roughly constant, $m_X \dot{N} \propto \frac{A \rho_0 r}{\tau_X}$.

•) We'll see that constraints are also quite insensitive to the details of $\rho(r)$, so we get similar results for sources with common M.

ii) Moving on to <u>annihilations</u>, let's consider



or :



χ^0 ... LSP neutralino

Tree-level, allowed by R-parity

Loop-level suppressed, DM does not couple to γ at tree-level.

•) For $m_X \gtrsim$ TeV (or for light final states like γ's) the $\langle \sigma v \rangle \simeq \frac{\alpha_{DM}^2}{m_X^2}$, which may be the <u>same</u> process that sets the thermal freeze-out abundance.

•) $$\frac{dN}{dt\,dV} \equiv \underbrace{\sigma}_{\text{x-section}} \underbrace{n_X v}_{\substack{\text{incident}\\ \text{flux}}} \underbrace{n_X}_{\substack{\text{density}\\ \text{of the target}}} \qquad \left(\text{for decays } \frac{n_X}{\tau_X} \right)$$

$$\Rightarrow \frac{dN_{ann}}{dt} = A \langle \sigma v \rangle \left(\frac{\rho_0}{m_X} \right)^2 (1\,\text{kpc})$$

$$\sim A\, 10^{-26} \frac{cm^3}{s} \left(\frac{0.4\,GeV}{m_X} \right)^2 (1\,\text{kpc})\, cm^{-6}$$

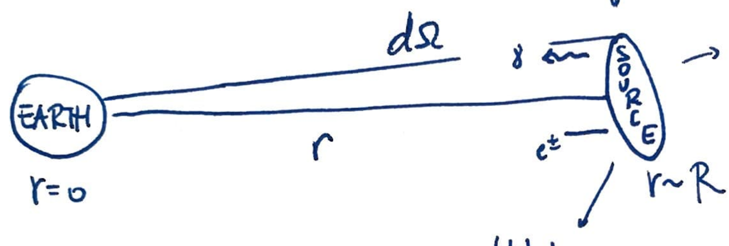
$$= 5 \cdot 10^{-8} s^{-1} \left(\frac{A}{m^2} \right) \left(\frac{\rho_0}{0.4\,GeV} \right)^2 \left(\frac{r}{\text{kpc}} \right) \left(\frac{TeV}{m_X} \right)^2 \left(\frac{\langle \sigma v \rangle}{10^{-26}\,cm^3 s^{-1}} \right).$$

$\sim$ 1 event / year.

J-FACTORS

•) A very common practice in the field is to factor out the astrophysical source dependence into a so-called J-factor.

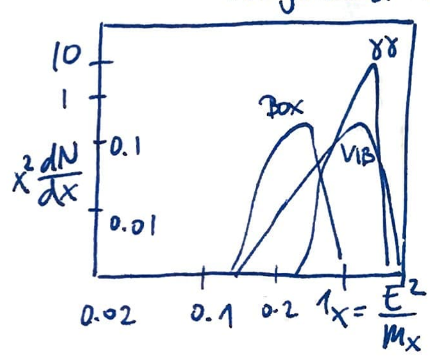
Here $XX \to \gamma\gamma, e^+e^-, WW, \ldots$



$\rightarrow n(\vec{r}) = \frac{\rho}{m_X}$

$\left.\frac{dN}{dE}\right|_s$

Bringmann & Weniger '12



•) The rate distribution is then:

$$\frac{dN}{dE\,dt\,dV} = \left.\frac{dN}{dE}\right|_s \left(\frac{A}{4\pi r^2}\right) \begin{cases} \langle\sigma v\rangle \frac{n_X^2}{2} & \text{, annihilations} \\ \frac{n_X}{\tau_X} & \text{, decays} \end{cases}$$

$$= \left.\frac{dN}{dE}\right|_s \left(\frac{A}{4\pi r^2}\right) \begin{cases} \frac{\langle\sigma v\rangle}{2m_X^2} \overbrace{\int_0^\infty \rho^2(r)\, dr\, d\Omega}^{8\pi J_{ann}} \\ \frac{1}{m_X \tau_X} \underbrace{\int_0^\infty \rho(r)\, dr\, d\Omega}_{\sim \frac{1}{R^2}\int \rho\, dV \sim \frac{M}{R^2}} \sim \frac{M}{m_X \tau_X R^2} \end{cases}$$

With NFW $\rho = \frac{\rho_0}{\left(\frac{r}{r_s}\right)\left(1+\frac{r}{r_s}\right)^2}$

•) We thus have for <u>annihilation</u>:

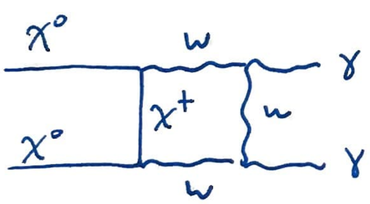
Dw Sph

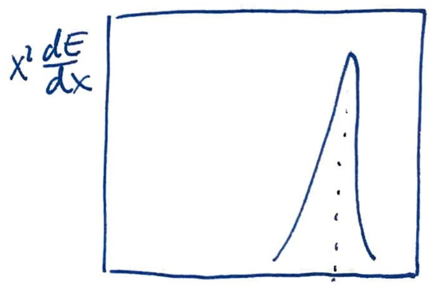
$$\frac{1}{A}\frac{dN_{ann}}{dE\,dt} = \frac{\langle\sigma v\rangle}{m_\chi^2}\left(\frac{dN}{dE}\right)_s J_{ann}\,, \qquad J_{ann} = \begin{cases} 10^{17-20}\ \frac{\mathrm{GeV}^2}{\mathrm{cm}^5} \\ 10^{22}\ \frac{\mathrm{GeV}^2}{\mathrm{cm}^5} \end{cases}$$

1° MW

The J-factors assumed the NFW profiles for dwarfs & MW center (1°). Local overdensities, like clumps with more DM may enhance by orders of mag.

•) Finally the energy spectra at the source $\left(\frac{dN}{dE}\right)_s$ depend on the microscopic process.

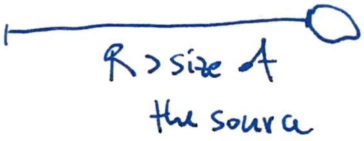




$$\left.\frac{dN}{dE}\right|_s \simeq 2\,\delta(E-m_\chi)$$

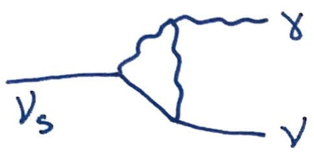
$x \sim 1 \qquad x = \frac{E}{m_\chi}$

•) For <u>decays</u>, the J-factor is simple:



$$\rho \sim \frac{M}{V}\,\delta(\rho - R) \qquad \int\!\!\int \rho\, dr\, d\Omega \sim \frac{1}{R^2}\int \rho\, dV = \frac{M}{R^2}$$

·) $\frac{1}{A}\frac{dN}{dE\,dt} = \frac{1}{4\pi}\frac{1}{m_x \tau_x}\left(\frac{dN}{dE}\right)_s \frac{M}{R^2}$

·) e.g. 

$\frac{dN}{dE} \sim \delta(E - m_x/2)$

*) For distant sources, the redshift $d_L(z)$ and γ absorption are relevant, too.

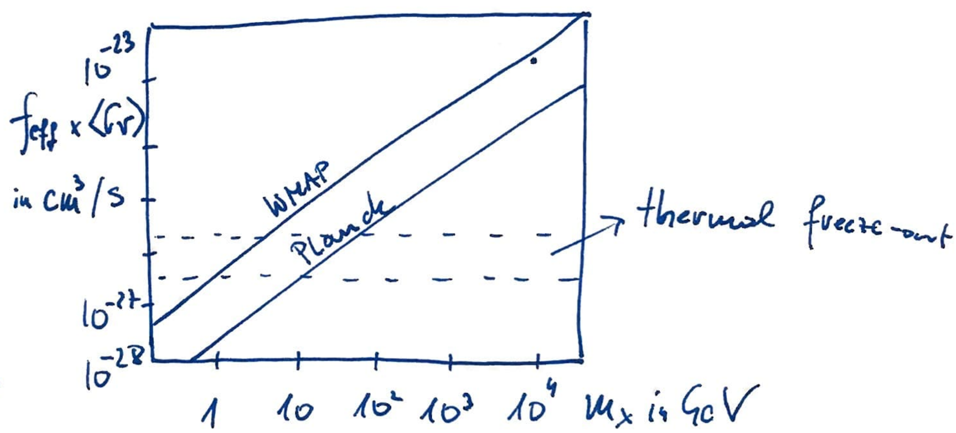
CURRENT LIMITS

·) The exact limit depends on the final state and the exposure of the experiment.

a) Early universe

WMAP & Planck bound from CMB effects.

@ $z_{ann} \sim 600$, $z_{dec} \sim 300$

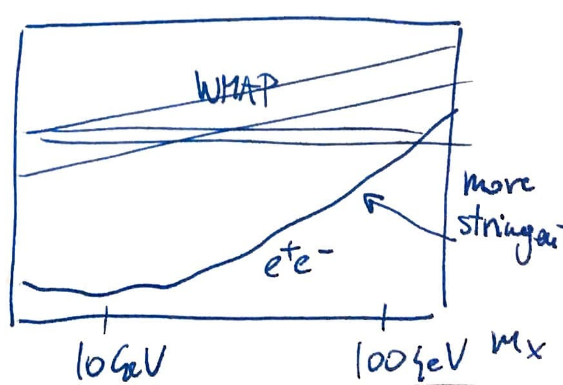


·) f_{eff} parametrizes the final state transmission to CMB

$$f_{eff} \cong \begin{cases} 0.4 & \text{for } e, W, top \\ 0.2 & \text{the rest, except } \gamma\text{'s} \\ \sim 0 & \text{for } \gamma\text{'s} \end{cases}$$

b) Charged leptons (galaxy propagation)

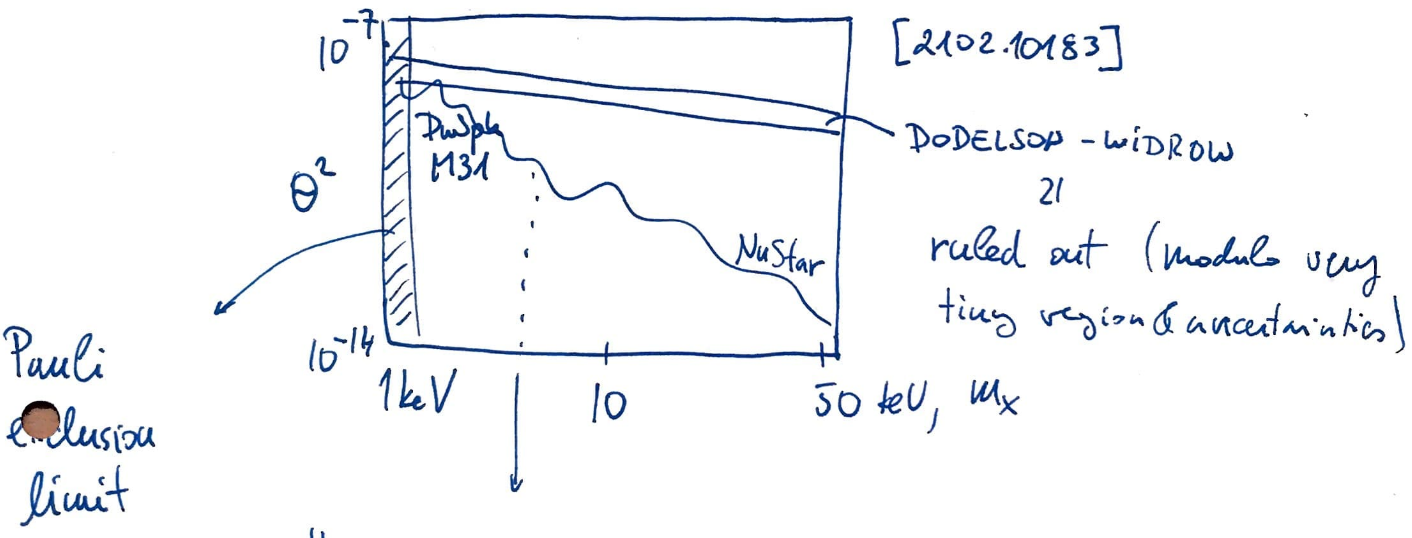
AMS '02 exp.



•) for the γ final state, the $\langle \sigma v \rangle$ is typically loop suppressed by at least $\frac{\alpha^2}{4\pi}$ (or more if there are heavy mediators in the loop)

FERMI : $\langle \sigma v \rangle_{\chi\chi \to \gamma\gamma} \lesssim 10^{-28} \frac{cm^3}{s}$ @ $m_\chi = 10\,GeV$

•) There are dedicated searches for X-rays, looking for photon monochromatic lines.



there was an exciting excess @ 3.5 keV, gone.

SUMMARY ON INDIRECT DETECTION

- Indirect detection experiments are looking for visible final state products of DM decay or annihilation.

- The rate of incoming particles depends on

$$\langle \sigma v \rangle n_\chi^2 \quad \text{for annihilation}$$

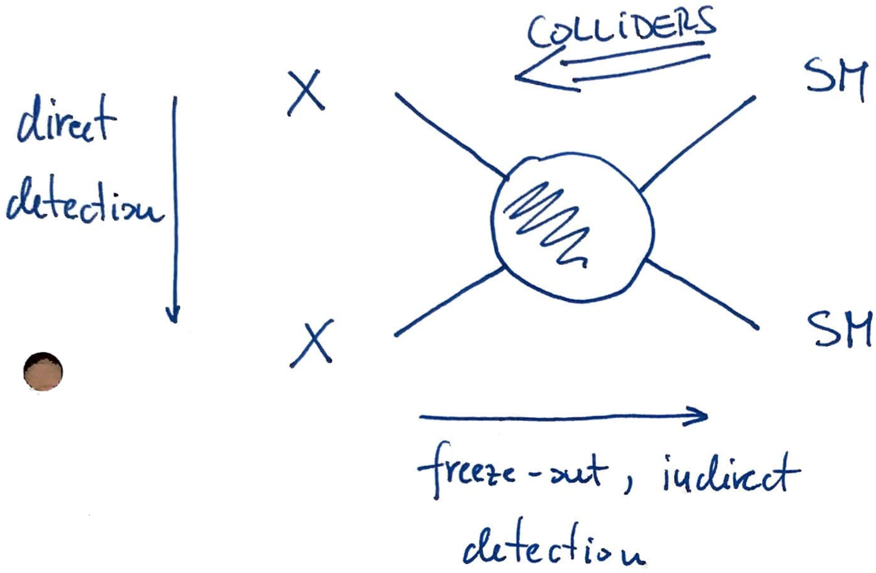
$$\frac{n_\chi}{\tau_\chi} \quad \text{for decays}$$

- The astrophysical uncertainties are stored in the J-factors, known for DwSph, MW 1°.

- Depending on the final state (like e^+e^-), the limits may be very stringent and even below the $\langle \sigma v \rangle$ for thermal freeze-out or below the Θ^2 for DW freeze-in.

Collider Searches for DM

•) We saw the 2-2 illustration for freeze-out, direct & indirect detection, the 3rd way is to collide SM-SM into X-X.

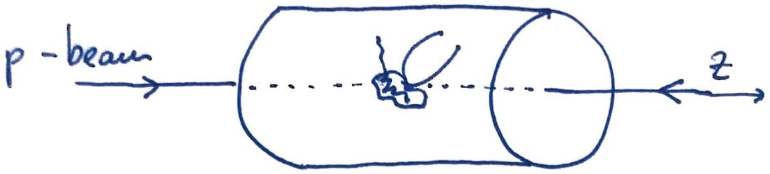


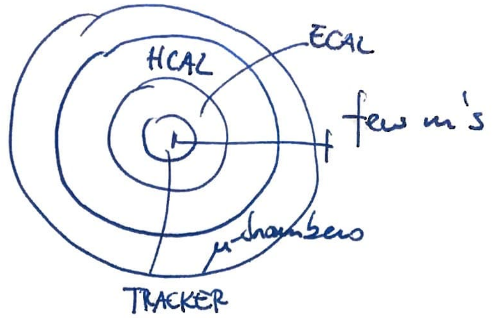
•) There is a range of "stable" particles we can collide, $e^\pm$, $p^\pm$ (even μ's).
LEP   LHC, TeVatron

•) For $m_X \in$ (GeV ~ TeV), the LHC is most suitable to detect since $\sqrt{s} = E_{CMS} \simeq 14$ TeV.

•) While DM may decay on τ_u scales, the typical collider lifetimes are $\sim \frac{\text{few meters}}{c} \ll \tau_u$, so DM if produced would be neutral (no charged tracks) and escapes detection, no ECAL/HCAL deposits.

•) Typical detector ATLAS, CMS





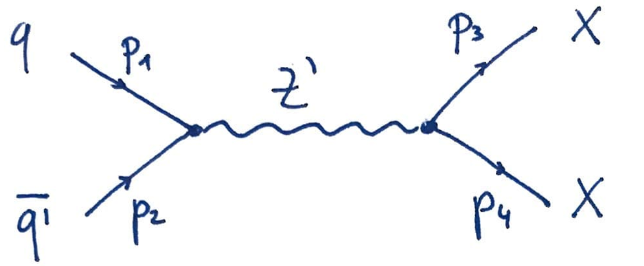
•) We know that E is transfered along the z axis, so the total ΣE in $\perp$ direction should vanish. The imbalance is called E_T^{miss} = missing transverse energy

•) A DM particle with $m_X \sim TeV$ or $E_X \sim TeV$ would take away a significant part of E and create a large imbalance $E_T^{miss} \sim m_X$ or E_X. While there are ν induced backgrounds in the SM, they are well known and predictable, $\Rightarrow$ we can look for S_{DM} / B_{SM}.

•) To perform a reliable collider study, one has to properly define a DM model and simulate S and B events. Then one devises cuts (typically large $\cancel{E}_T$) to enhance $\frac{S}{B}$. We will do simple $\sigma \times \mathcal{L} = N_{events}$ estimates instead.

SIMPLE COLLIDER ESTIMATE

•) Consider a simple toy model with a Z' & X=DM.



$$\hat{s} = (p_1 + p_2)^2$$

Let's define the coupling to be vectorial with

g_1 to quarks and g_2 for X

$$j_q^{Z'\mu} = g_1 \bar{q} \gamma^\mu q \quad , \quad j_X^{Z'\mu} = g_2 \bar{X} \gamma^\mu X$$

Since we do not know $M_{Z'}$, let's consider its full mass dependence via a Breit-Wigner (BW)

$$\text{•) Partonic } \hat{\sigma}_{qq' \to XX} = \underbrace{3\pi \left(\frac{g_1^2}{4\pi}\right)}_{\text{SM-SM}} \underbrace{\frac{\Gamma M}{(\hat{s} - M^2)^2 + (\Gamma M)^2}}_{\text{MEDIATOR}} \underbrace{Br(Z' \to XX)}_{\text{DM XX part}}$$

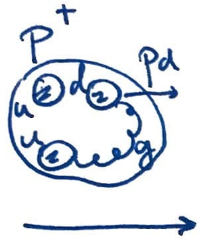
•) The BW sets the dimension : $\begin{cases} \sim \frac{1}{\hat{s}} & \text{for massless like } \gamma \\ \sim \frac{1}{M^2} & \text{for massive, } \hat{s} < M^2 \end{cases}$

•) The DM part is in the

$$Br_{Z' \to XX} = \frac{1}{\Gamma_{tot}} \frac{g_2^2}{48\pi} M \underbrace{\left(1 - \frac{m_X^2}{\hat{s}}\right)^2}_{\text{phase space suppression}} \left(1 + \frac{m_X^2}{2\hat{s}}\right).$$

$\Gamma_{tot} \propto M$ so $[Br] = 0$

b) Now that we have the partonic $\hat{\sigma}_{qq \to xx}$, we can convolute it with the Parton Distribution Functions PDFs of the proton $f_q(x, Q^2)$, which give the final $\sigma(pp \to XX)$ at the nucleon level.

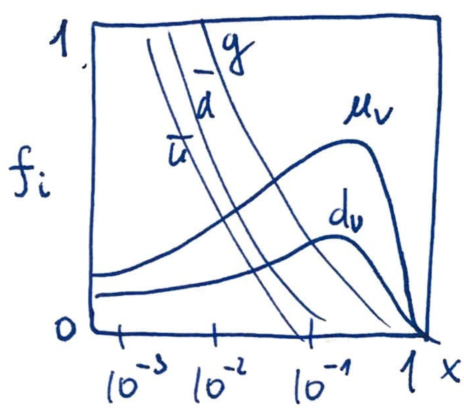


$p_i = x_i P$ ... partonic momentum fraction

p_i ... parton momentum

P... momentum of the proton $\sqrt{s} = (P_1 + P_2)^2$

PROTON PDFS [PDG '21]



*) light valence f_q peak at $x \sim 0.1$

*) $\int f_u \sim 2 \int f_d$

•) Putting it all together, we have

$$\sigma_{pp \to xx} = \frac{1}{9} \int_{\frac{m_x^2}{s}}^{1} dx_1 \int_{\frac{m_x^2}{x_1 s}}^{1} dx_2 \sum_{q, q'} \hat{\sigma}_{qq' \to xx} \left(f_q(x_1) f_{\bar{q}}(x_2) + (1 \leftrightarrow 2) \right)$$

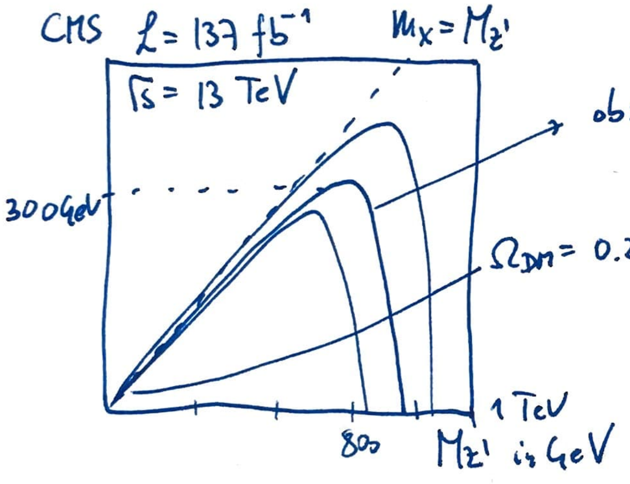
•) That's it! From $\mathcal{L} \times \sigma_{pp \to xx} = N_{events}$

$\mathcal{L}$ ... luminosity measured at the LHC, $\mathcal{L} \gtrsim 100 fb^{-1}$ $\propto$ amount of statistics collected.

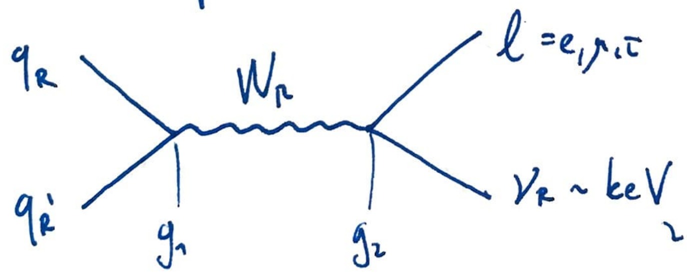
•) Finally, we have to devise cuts to enhance the signal rate (say $\not{E}_T \gtrsim 100\,\text{GeV}$, initial jets, γ's, W's), which typically reduces the rate by $\varepsilon \sim 1\% - 10\%$.

$$N_{obs} = \varepsilon \cdot \mathcal{L} \cdot \sigma$$

Current searches exist in a myriad of final states, here we show [CMS-PAS-EXO-19-003]



•) Another example is ν_R as DM in the LRSM



$g_1 = g_2 = g_R \sim 0.65$

$\sim m_X \sim 0$

$$M_{W_R} > \begin{cases} 4.6\,\text{TeV} & \mu^- \\ 4.9\,\text{TeV} & e^- \end{cases}$$

[1801.05813]