MAX-PLANCK-INSTITUT FÜR KERNPHYSIK

Neutrinos and Abelian Gauge Symmetries

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Symmetries of the Standard Model

The Standard Model (SM), based on the gauge symmetry group $G_{\rm SM} \equiv SU(3)_C \times SU(2)_L \times U(1)_Y$, features an accidental abelian global symmetry $U(1)^4$, connected to the conservation of baryon number B and individual lepton numbers $L_{e,\mu,\tau}$. Since $B+L$ is broken by quantum effects, the abelian symmetry group of the SM takes the form

$$\mathcal{G} \equiv U(1)_{B-L} \times U(1)_{L_e-L_\mu} \times U(1)_{L_\mu-L_\tau} .$$

Following the success of the gauge principle in particle physics, we can try to promote this *global* symmetry to a *local* symmetry; anomaly cancellation then requires the introduction of three right-handed neutrinos ν_R, which render the full gauge group $G_{\rm SM} \times \mathcal{G}$ anomaly free.[1] Furthermore, the right-handed neutrinos automatically lead to non-zero neutrino masses, alleviating one of the main shortcomings of the SM:

$$m_\nu \neq 0 .$$

The connection goes further: the unflavored gauge symmetry $U(1)_{B-L}$ is linked to the nature of neutrinos, i.e. whether they are Majorana ($\nu = \overline{\nu}$) particles with $\Delta(B-L) = 2$ or Dirac ($\nu \neq \overline{\nu}$) particles with $\Delta(B-L) = 0, 4, 6, \dots$ All cases are closely connected to the matter–antimatter asymmetry of our Universe.[2,3]

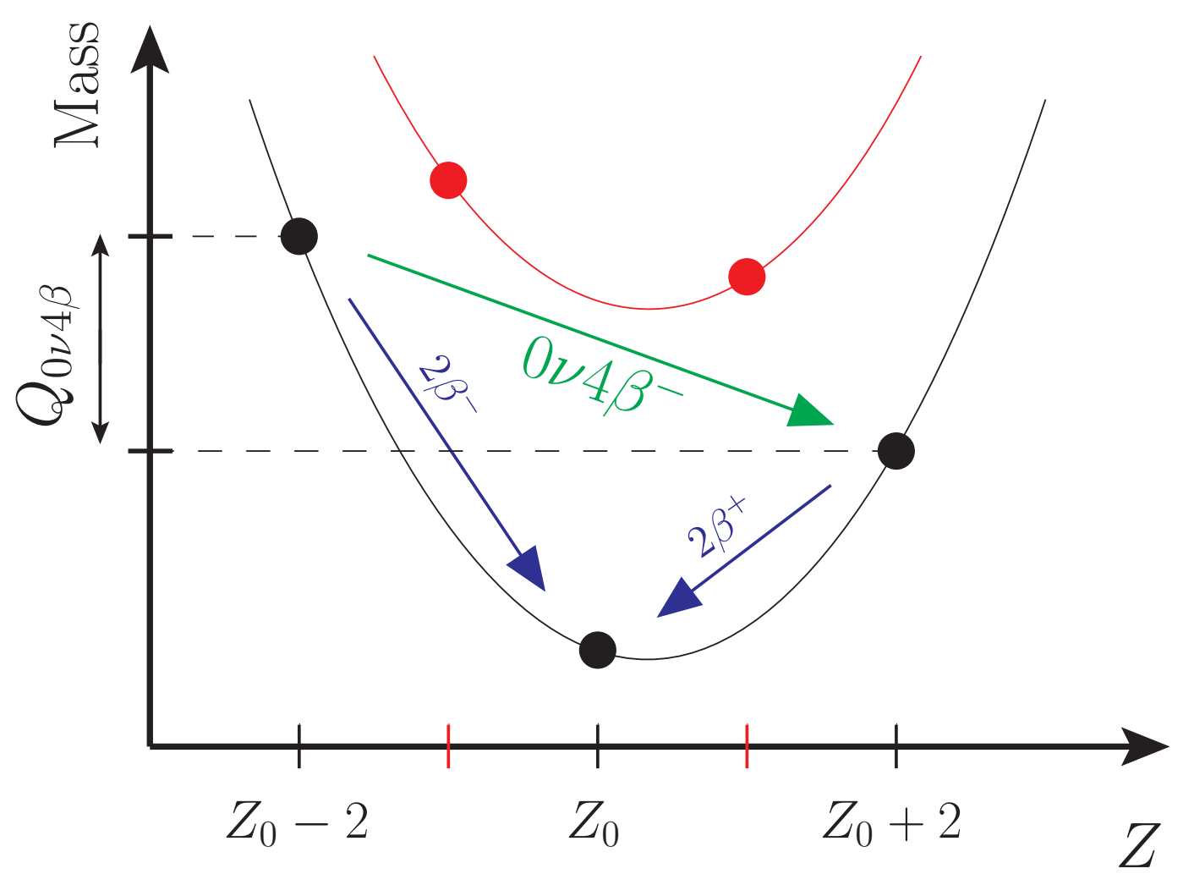
Flavored subgroups $U(1)' \subset \mathcal{G}$ can shed light on the peculiar leptonic mixing pattern and mass ordering in a simple and testable manner, or even enforce viable texture zeros in the neutrino mass matrices.[1,4,5]

Beyond $\mathcal{G}$, abelian symmetries $U(1)_{\rm DM}$ in the dark matter sector can give rise to naturally light sterile neutrinos, which provide a new portal between visible and dark sector, and also resolve some longstanding anomalies in neutrino experiments.[7]

Unflavored Gauge Symmetries

The symmetry $U(1)_{B-L}$ is connected to the Dirac vs. Majorana nature of neutrinos. If broken by *two* units, $\Delta(B-L) = 2$, one obtains Majorana neutrinos and the signature decay of neutrinoless double beta decay $0\nu2\beta$. If, on the other hand, neutrinos are Dirac fermions, $B-L$ is either unbroken, or broken by 4, 6, 8, …, $\neq 2$ units, still allowing for lepton number violating interactions.

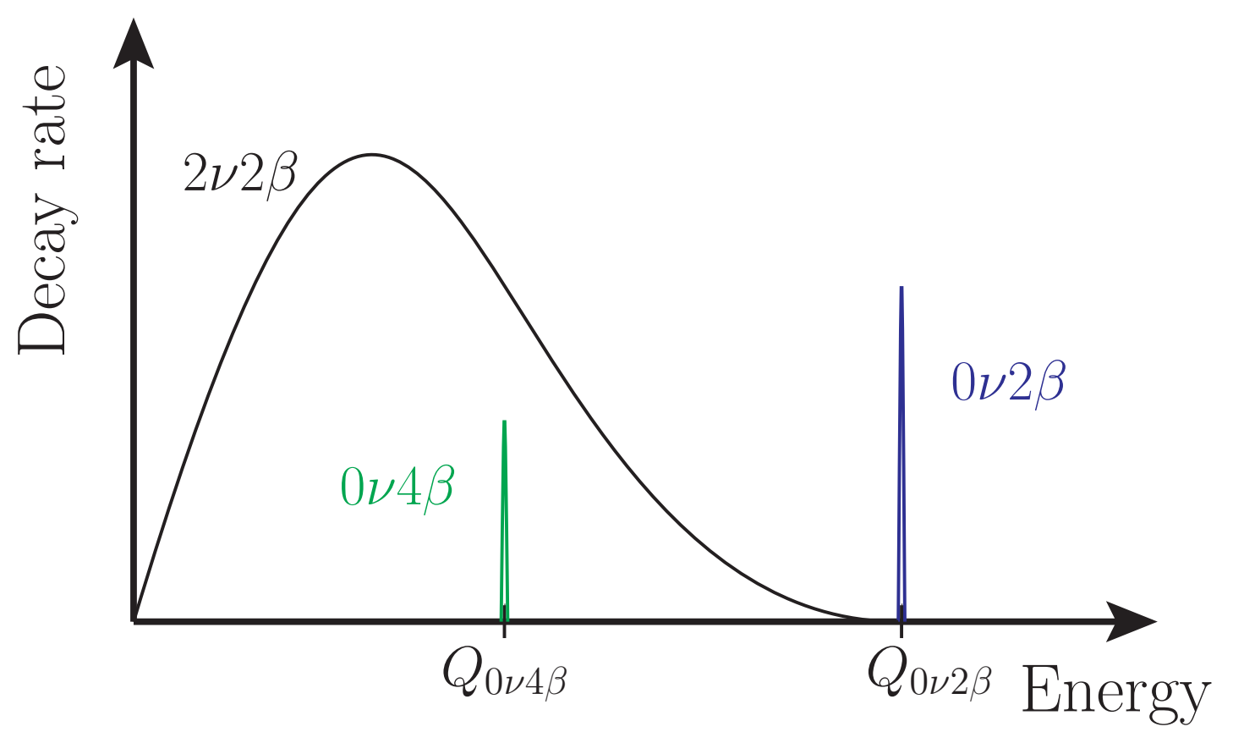
A UV-complete realization of such lepton-number-violating Dirac neutrinos can be readily constructed, and in particular give rise to a new Dirac leptogenesis mechanism.[2] At low energies, lepton number violation $\Delta L = 4$ is difficult to test, but in principle there exists the analogue to $0\nu2\beta$: neutrinoless *quadruple* beta decay $0\nu4\beta$.



Only three isotopes (A,Z) can undergo this $\Delta L = 4$ decay

$$(A,Z) \to (A, Z+4) + 4e^- ,$$

and four more offer electron-capture modes (see table to the right). The experimental signature of $0\nu4\beta$ are four emitted electrons, whose energies sum up to the Q value and result in a line. Without tracking, the standard $2\nu2\beta$ decay will overlap this energy region, but first upper limits on this novel $\Delta L = 4$ decay can be obtained with existing data.[3]



Lepton number violation $\Delta L \neq 2$ is not only difficult to test experimentally, but the simplest theoretical realizations lead to unobservably small rates. Nonetheless, the quest for lepton number violation need not be over once we discover neutrinos to be Dirac particles.

element	$Q_{0\nu4\beta}$	other decays	NA in %
${}^{96}_{40}$Zr → ${}^{96}_{44}$Ru	0.629 MeV	$\tau^{2\nu2\beta}_{1/2} \simeq 2\times10^{19}$ yr	2.8
${}^{136}_{54}$Xe → ${}^{136}_{58}$Ce	0.044 MeV	$\tau^{2\nu2\beta}_{1/2} \simeq 2\times10^{21}$ yr	8.9
${}^{150}_{60}$Nd → ${}^{150}_{64}$Gd	2.079 MeV	$\tau^{2\nu2\beta}_{1/2} \simeq 7\times10^{18}$ yr	5.6
	$Q_{0\nu4\mathrm{EC}}$		
${}^{124}_{54}$Xe → ${}^{124}_{50}$Sn	0.577 MeV	–	0.095
${}^{130}_{56}$Ba → ${}^{130}_{52}$Te	0.090 MeV	$\tau^{2\nu2\mathrm{EC}}_{1/2} \sim 10^{21}$ yr	0.106
${}^{148}_{64}$Gd → ${}^{148}_{60}$Nd	1.138 MeV	$\tau^{\alpha}_{1/2} \simeq 75$ yr	–
${}^{154}_{66}$Dy → ${}^{154}_{62}$Sm	2.063 MeV	$\tau^{\alpha}_{1/2} \simeq 3\times10^{6}$ yr	–
	$Q_{0\nu3\mathrm{EC}\beta^+}$		
${}^{148}_{64}$Gd → ${}^{148}_{60}$Nd	0.116 MeV	$\tau^{\alpha}_{1/2} \simeq 75$ yr	–
${}^{154}_{66}$Dy → ${}^{154}_{62}$Sm	1.041 MeV	$\tau^{\alpha}_{1/2} \simeq 3\times10^{6}$ yr	–
	$Q_{0\nu2\mathrm{EC}2\beta^+}$		
${}^{154}_{66}$Dy → ${}^{154}_{62}$Sm	0.019 MeV	$\tau^{\alpha}_{1/2} \simeq 3\times10^{6}$ yr	–

Table: Candidates for the nuclear $\Delta L = 4$ processes neutrinoless quadruple beta decay and electron capture, the corresponding Q values, competing (observed) decay channels with half-life $\tau^j_{1/2}$, and natural abundance (NA) of the candidate isotopes.

Flavored Gauge Symmetries

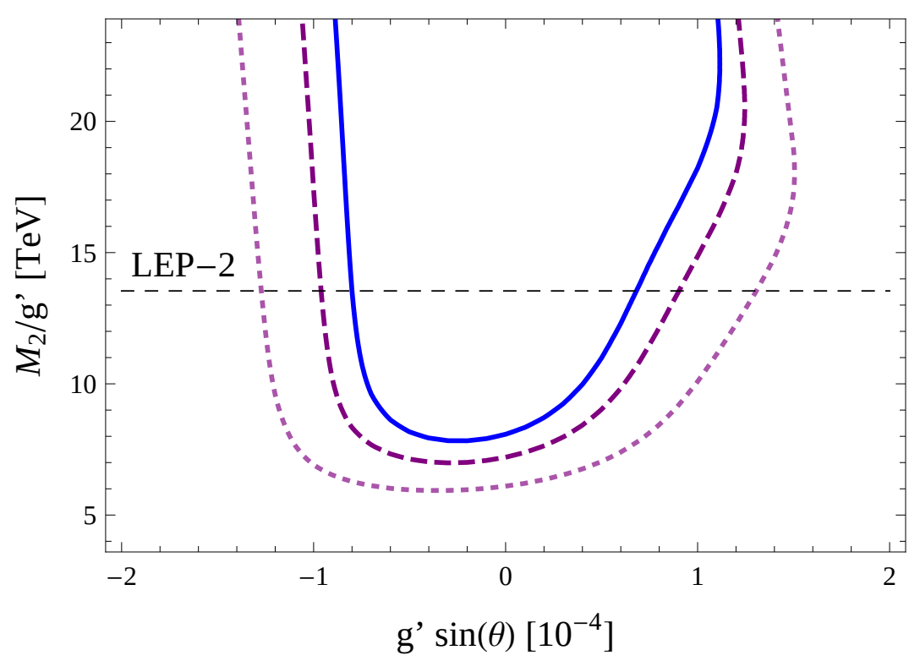
As an alternative to *discrete non-abelian global* symmetries one can employ *continuous abelian local* symmetries $U(1)' \subset \mathcal{G}$ to shed light on the peculiar leptonic mixing pattern in a very simple and testable manner. Such approaches are typically less predictive than abelian flavor symmetry models, but come with much smaller parameter count and particle content. Each neutrino hierarchy exhibits its own approximately conserved lepton charge.

Normal Ordering

For positive Δm^2_{32} one finds an approximately L_e symmetric Majorana mass matrix

$$m_\nu^{L_e} \simeq \begin{pmatrix} 0 & 0 & 0 \\ 0 & \times & \times \\ 0 & \times & \times \end{pmatrix} ,$$

which leads to small $0\nu2\beta$ rates, induced by $\mathcal{O}(10^{-2})$ corrections to m_ν. This structure can be enforced by an anomaly-free gauge symmetry $U(1)_{B+3(L_e-L_\mu-L_\tau)} \subset \mathcal{G}$, spontaneously broken by an SM-singlet scalar S; the seesaw mechanism gives an approximate $m_\nu^{L_e}$ after the $\langle S\rangle$-induced correction ΔM_R to $M_R^{L_e-L_\mu-L_\tau}$. The mass M_2 of the new gauge boson Z' and its mixing angle θ with the Z are constrained by direct searches and electroweak precision data.[4]

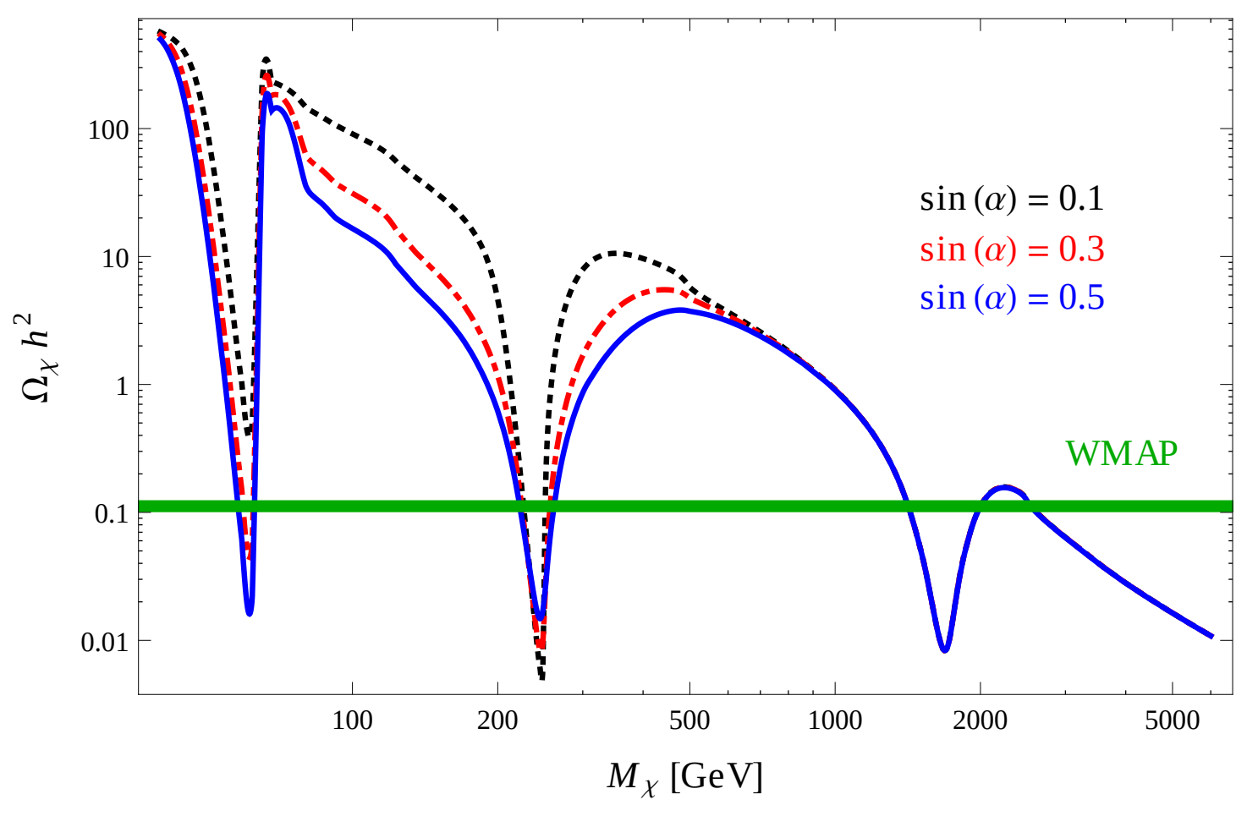


Inverted Ordering

A negative Δm^2_{32} yields the distinct pattern

$$m_\nu^{L_e-L_\mu-L_\tau} \simeq \begin{pmatrix} 0 & \times & \times \\ \times & 0 & 0 \\ \times & 0 & 0 \end{pmatrix} .$$

This most difficult pattern can be realized again with a gauged $U(1)_{B+3(L_e-L_\mu-L_\tau)}$, but extended by two more right-handed neutrinos, one decoupled by a $\mathbb{Z}_2$. The decoupled state χ forms Majorana dark matter and the correct relic density can be obtained around either of the scalar resonances (α denoting the h–s scalar mixing angle) or the Z' resonance.[4]

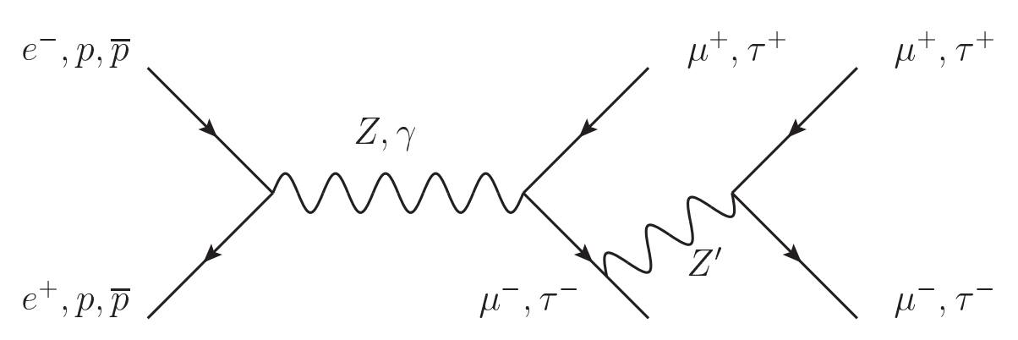


Quasi-Degenerate

For neutrino masses $m_i \gg 0.1$ eV there emerges yet another symmetry:

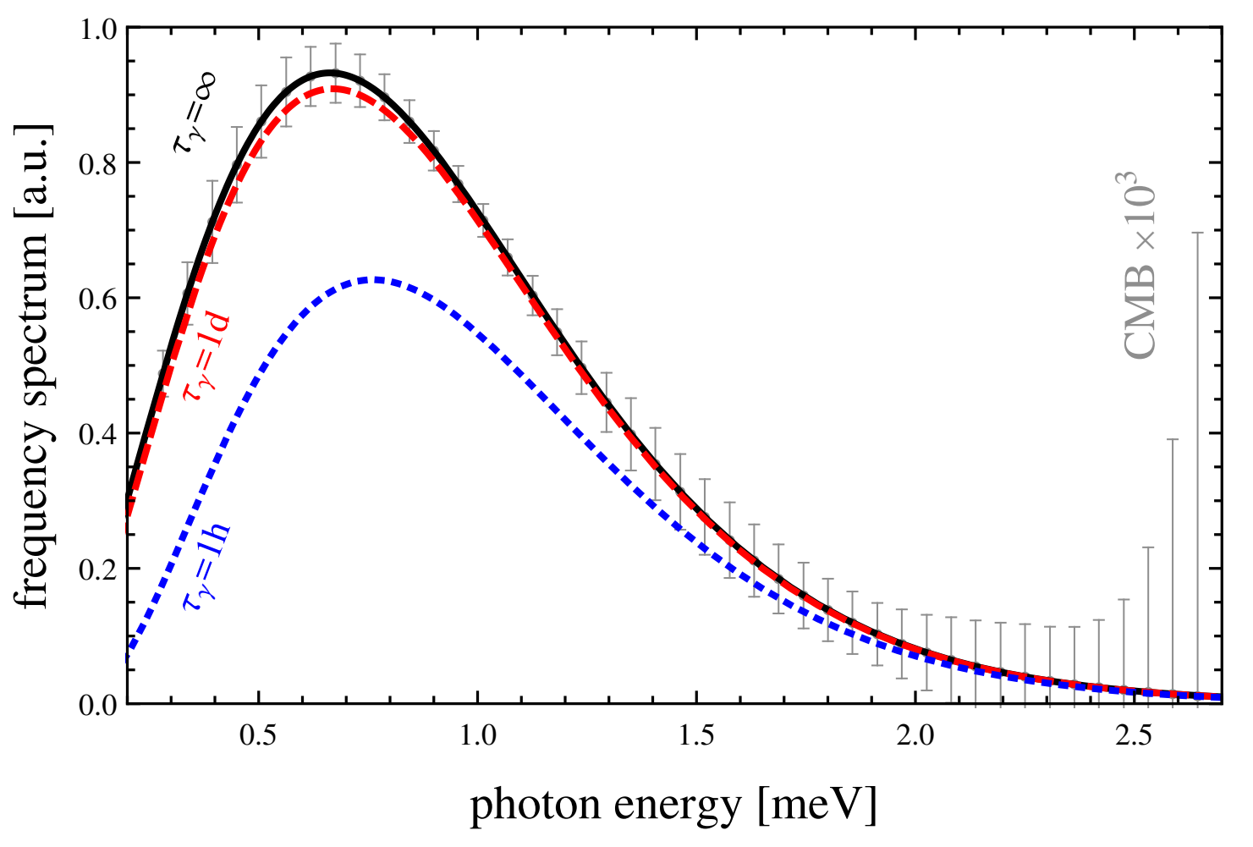
$$m_\nu^{L_\mu-L_\tau} \simeq \begin{pmatrix} \times & 0 & 0 \\ 0 & 0 & \times \\ 0 & \times & 0 \end{pmatrix} ,$$

which leads to large $0\nu2\beta$ rates and close-to-maximal θ_{23}. The Z' boson of the associated anomaly-free gauge symmetry $U(1)_{L_\mu-L_\tau}$ contributes to the muon's anomalous magnetic moment and can resolve the long-standing theory–experiment discrepancy for values $M_{Z'}/g' \simeq 200$ GeV, and is furthermore testable at the LHC by means of final-state muon radiation.[5]



Lifetime of Light

Massive photons form a simple extension of the Standard Model, theoretically easily achievable and experimentally constrained to small values, namely $m < 10^{-18}\,\mathrm{eV} \simeq 2\times10^{-54}$ kg from measurements of the solar magnetic field. With a nonzero mass, even of such tiny proportions, photons are no longer necessarily *stable*, and are kinematically allowed to decay, e.g. into two of the lightest neutrinos.



A model-independent bound on the photon's lifetime τ_γ can be obtained by studying deviations of the well-measured cosmic microwave background from the usual black-body spectrum, resulting in

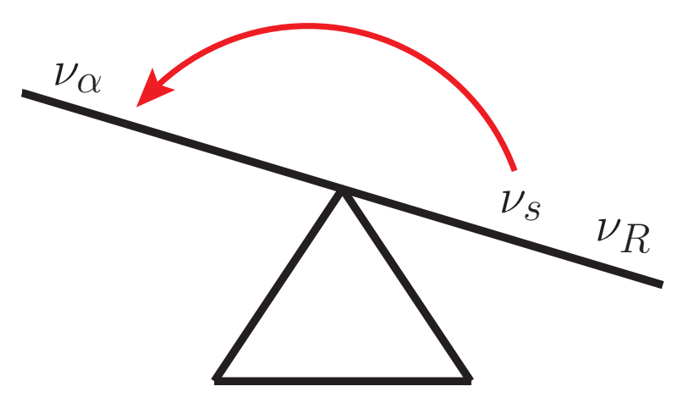
$$\tau_\gamma > 3\,\mathrm{yr} \left(\frac{m}{10^{-18}\,\mathrm{eV}}\right) .$$

This short lifetime only holds in the photon's rest frame and increases by a factor 10^{18} for photons in the visible spectrum.[6]

Light Sterile Neutrinos and Dark Symmetries

A number of short-distance experiments seem incompatible with the standard three-neutrino paradigm and hint at the existence of one or more *sterile* neutrinos with masses around 1 eV and $\mathcal{O}(0.1)$ mixing with the active neutrinos. The lightness of these steriles then begs for a theoretical explanation.

Seeing as small active-neutrino masses are beautifully explained by the seesaw mechanism, it is only natural to use that same seesaw to suppress the sterile neutrino mass, i.e. put a ν_s on the light side of the seesaw.



This *minimal extended seesaw* requires a particular structure in the $(\nu_L, \nu_R^c, \nu_S^c)$ mass matrix

$$M_{\rm MES} = \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & M_R & m_S \\ 0 & m_S^T & 0 \end{pmatrix} ,$$

and yields seesaw-suppressed active (sterile) neutrino masses $m_\nu = \mathcal{O}(m_D^2/M_R)$ ($m_s = \mathcal{O}(m_S^2/M_R)$) and active–sterile mixing $\mathcal{O}(m_D/m_S)$. The MES structure can be realized using simple $U(1)'$ symmetries at the TeV scale and automatically gives the required $\mathcal{O}(0.1)$ active–sterile mixing.

Anomaly cancellation requires additional fermions, which form automatically stable multicomponent dark matter, interacting with the SM via the kinetic-mixing, Higgs, and sterile-neutrino portal.[7]

References

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