Direct detection of fourth generation Majorana neutrino dark matter

Yu-Feng Zhou[*]

State Key Laboratory of Theoretical Physics,
Kavli Institute for Theoretical Physics China,
Institute of Theoretical Physics, Chinese Academy of Sciences,
Beijing, 100190, P.R. China
E-mail: yfzhou@itp.ac.cn

Heavy stable fourth generation Majorana neutrinos contribute to a small fraction of the relic density of dark matter (DM) in the Universe. Due to its relatively strong coupling to the standard model particles, it can be probed by the current direct and indirect DM detection experiments even it is a subdominant component of the halo DM. We show that the current Xenon100 data constrain the mass of the stable Majorana neutrino to be greater than the mass of the top quark. The effective spin-independent cross section for the neutrino elastic scattering off nucleon is predicted to be $\sim 1.5 \times 10^{-44}$ cm^2, which is insensitive to the neutrino mass and mixing and can be reached by the direct DM detection experiments in the near future. In the same mass region the predicted effective spin-dependent cross section for the heavy neutrino scattering off proton is in the range of 2×10^{-40} cm$^2 \sim 2 \times 10^{-39}$ cm^2, which is within the reach of the ongoing DM indirect search experiments. We demonstrate such properties of the heavy neutrino DM in a fourth generation model with the stability of the fourth Majorana neutrino protected by an anomaly-free $U(1)$ gauge symmetry.

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[*]Speaker.

Models with chiral fourth generation fermions are among the simplest and well-motivated extensions of the standard model (SM) and have been extensively studied [1]. The condition for CP symmetry violation in the SM requires at least three generations of fermions [2]. But, there is no upper limit on the number of generations from the first principle. In the SM the amount of CP violation is not large enough to explain the baryon-antibaryon asymmetry in the Universe. The inclusion of fourth generation quarks leads to two extra CP phases in quark sector and possible larger CP violation [3], which is helpful for electroweak baryogenesis. With very massive quarks in the fourth generation, it has been proposed that the electroweak symmetry breaking may become a dynamical feature of the SM [4–7].

Heavy stable neutrinos with mass greater than ~ 1 GeV are possible candidates for the cold DM [8, 9]. However, if the neutrino is the dominant component of the halo DM, the current DM direct search experiments have imposed strong constraints on its mass [10–15]. On the other hand, it is well-known that for a neutrino heavier than $\sim m_Z/2$, the cross section for its annihilation is in general too large to reproduce the observed DM relic density. For the neutrino heavier than m_W, the contribution from $f\bar{f}$ channels decrease rapidly, but other channels such as $W^{\pm}W^{\mp}$, $Z^0 h^0$ etc. are opened. For these processes the corresponding cross section does not decrease with the increasing of the neutrino mass, resulting in a relic density always decreases with the growing of the neutrino mass, and a thermal relic density far below the observed total DM relic density [16]. Thus the neutrino DM can only contribute to a small fraction of the relic density of DM and a small fraction of the halo DM density as well. Despite its very low number density in the halo, it can still be probed by the underground DM direct detection experiments due to its relatively strong coupling to the target nuclei, which provides a way to search for new physics beyond the SM complementary to the LHC.

In this talk, we discuss the consequence of this possibility in a model with a fourth generation Majorana neutrino DM. The stability of the fourth Majorana neutrino protected by an additional generation-dependent $U(1)$ gauge symmetry which is anomaly-free. The details of the analysis can be found in Ref. [17]. We consider a simple extension of the SM with a sequential fourth generation and an additional $U(1)_F$ gauge symmetry. The $U(1)$ extensions to the SM are well motivated from the point view of grand unification such as the $SO(10)$ and E_6 and have rich phenomenology [18] which can be reached by the on going LHC experiments. The flavor contents in the model are given by

$$q_{iL} = \begin{pmatrix} u_{iL} \\ d_{iL} \end{pmatrix}, \; \ell_{iL} = \begin{pmatrix} v_{iL} \\ e_{iL} \end{pmatrix}, \; u_{iR}, \; d_{iR}, \; v_{iR}, \; e_{iR} \; (i = 1, \ldots, 4). \tag{1}$$

All the fermions in the model are vector-like under the extra gauge interactions associated with $U(1)_F$. The $U(1)_F$ charges of the fermions could be generation-dependent. In order to evade the stringent constraints from the tree-level flavor changing neutral currents (FCNCs), the $U(1)_F$ charges Q_{qi} for the first three generation quarks are set to be the same, i.e. $Q_{qi} = Q_q$, $(i = 1, 2, 3)$ while $Q_{q4} = -3Q_q$ for the fourth generation quarks. Similarly, the $U(1)_F$ charges for the first three generation and the fourth generation leptons are Q_L and $-3Q_L$, respectively. In general, Q_q and Q_L can be different. For simplicity, we consider $Q_q = Q_L = 1$. With this set of flavor contents and $U(1)_F$ charge assignments, it is straight forward to see that the new gauge interactions are anomaly-free. Since the gauge interaction of $U(1)_F$ is vector-like, the triangle anomalies of

$[U(1)_F]^3$, $[SU(3)_C]^2 U(1)_F$ and $[\text{gravity}]^2 U(1)_F$ are all vanishing. The anomaly of $U(1)_Y[U(1)_F]^2$ is zero because the $U(1)_Y$ hypercharges cancel for quarks and leptons separately in each generation, namely $\sum(-Y_{qL} + Y_{qR}) = 0$ and $\sum(-Y_{\ell L} + Y_{\ell R}) = 0$. The anomaly of $[SU(2)_L]^2 U(1)_F$ is also zero due to the relation $\sum_{i=1}^{4} Q_{qi} = 0$ and $\sum_{i=1}^{4} Q_{Li} = 0$. Thus in this model, the gauge anomalies generated by the first three generation fermions are canceled by that of the fourth generation one, which also gives a motivation for the inclusion of the fourth generation.

The gauge symmetry $U(1)_F$ is to be spontaneously broken by the Higgs mechanism. For this purpose we introduce two SM singlet scalar fields $\phi_{a,b}$ which carry the $U(1)_F$ charges $Q_a = -2Q_L$ and $Q_b = 6Q_L$ respectively. The $U(1)_F$ charges of $\phi_{a,b}$ are arranged such that ϕ_a can have Majorana type of Yukawa couplings to the right-handed neutrinos of the first three generations ν_{iR} ($i = 1, 2, 3$) while ϕ_b only couples to the fourth generation neutrino ν_{4R}. After the spontaneous symmetry breaking, the two scalar fields obtain vacuum expectation values (VEVs) $\langle \phi_{a,b} \rangle = v_{a,b}/\sqrt{2}$. The relevant interactions in the model are given by

$$\begin{aligned}
\mathscr{L} &= \bar{f}_i i \gamma^\mu D_\mu f_i + (D_\mu \phi_a)^\dagger (D_\mu \phi_a) + (D_\mu \phi_b)^\dagger (D_\mu \phi_b) \\
&- Y_{ij}^d \bar{q}_{iL} H d_{iR} - Y_{ij}^u \bar{q}_{iL} \tilde{H} u_{iR} - Y_{ij}^e \bar{\ell}_{iL} H e_{iR} - Y_{ij}^\nu \bar{\ell}_{iL} \tilde{H} \nu_{iR} \\
&- \frac{1}{2} Y_{ij}^m \overline{\nu_{iR}^c} \phi_a \nu_{jR} \, (i, j = 1, 2, 3) - \frac{1}{2} Y_4^m \overline{\nu_{4R}^c} \phi_b \nu_{4R} - V(\phi_a, \phi_b, H) + \text{H.c.}
\end{aligned} \tag{2}$$

where f_i stand for left- and right-handed fermions, and H is the SM Higgs doublet. $D_\mu f_i = (\partial_\mu - ig_1 \tau^a W_\mu^a - iY g_2 B_\mu - iQ_f g_F Z_\mu') f_i$ is the covariant derivative with Z_μ' the extra gauge boson associated with the $U(1)_F$ gauge symmetry, and g_F the corresponding gauge coupling constant. Since $\phi_{a,b}$ are SM singlets, they do not play any role in the electroweak symmetry breaking. Thus Z' obtains mass only from the VEVs of the scalars

$$m_{Z'}^2 = g_F^2 (Q_a^2 v_a^2 + Q_b^2 v_b^2). \tag{3}$$

From the $U(1)_F$ charge assignments in the model, the four by four Yukawa coupling matrix is constrained to be of the block diagonal form $\mathbf{3} \otimes \mathbf{1}$ in the generation space. Since the $U(1)_F$ charges are the same for the fermions in the first three generation, there is no tree level FCNC induced by the Z'-exchange in the physical basis after diagonalization. Thus a number of constraints from the low energy flavor physics such as the neutral meson mixings and the $b \to s\gamma$ can be avoided.

The direct search for the process $e^+ e^- \to Z' \to \ell^+ \ell^-$ at the LEP-II leads to a lower bound on the ratio of the mass to the coupling to leptons: $M_{Z'}/g_F \geq 6$ TeV [19] for vector-like interactions, which corresponds to a more stringent lower bound: $\sqrt{Q_a^2 v_a^2 + Q_b^2 v_b^2} \geq 6$ TeV. The current searches for narrow resonances in the Drell-Yan process $pp \to Z' \to \ell^+ \ell^-$ at the LHC impose an alternative bound on the mass and the couplings of the Z' boson. For a model with sequential neutral gauge boson Z'_{SSM} which by definition has the same couplings as that for the SM Z^0 boson [18], the latest lower bounds on its mass $M_{Z'_{SSM}}$ is 1.94 TeV from CMS [20] and 1.83 TeV from ATLAS [21] respectively. The bound on $M_{Z'_{SSM}}$ can be translated into the bound on the mass and couplings of the Z' in this model. For instance, $g_F \lesssim 0.029$ for $M_{Z'} = 1.44$ TeV and $g_F \lesssim 0.0051$ for $M_{Z'} = 0.94$ TeV, respectively.

The fourth generation neutrinos obtain both Dirac and Majorana mass terms through the vacuum expectation values (VEVs) of H and ϕ_b. In the basis of $(\nu_L, \nu_R^c)^T$ the mass matrix for the

fourth neutrino is given by

$$m_v = \begin{pmatrix} 0 & m_D \\ m_D & m_M \end{pmatrix}, \qquad (4)$$

where $m_D = Y_4^v v_H / \sqrt{2}$ with $v_H = 246$ GeV and $m_M = Y_4^m v_{\phi_b} / \sqrt{2}$. The left-handed components $(v_{1L}^{(m)}, v_{2L}^{(m)})$ of the two mass eigenstates are related to the ones in the flavor eigenstates by a rotation angle θ

$$v_{1L}^{(m)} = -i(c_\theta v_L - s_\theta v_R^c), \quad v_{2L}^{(m)} = s_\theta v_L + c_\theta v_R^c, \qquad (5)$$

where $s_\theta \equiv \sin\theta$ and $c_\theta \equiv \cos\theta$. The value of θ is defined in the range $(0, \pi/4)$ and is determined by

$$\tan 2\theta = \frac{2m_D}{m_M}, \qquad (6)$$

with $\theta = 0$ $(\pi/4)$ corresponding to the limit of minimal (maximal) mixing. The phase i is introduced to render the two mass eigenvalues real and positive. The two Majorana mass eigenstates are $\chi_1 = v_{1L}^{(m)} + v_{1L}^{(m)c}$ and $\chi_2 = v_{2L}^{(m)} + v_{2L}^{(m)c}$, respectively. The masses of the two neutrinos are given by $m_{1,2} = (\sqrt{m_M^2 + 4m_D^2} \mp m_M)/2$. In terms of the mixing angle θ they can be rewritten as $m_1 = (s_\theta/c_\theta)m_D$ and $m_2 = (c_\theta/s_\theta)m_D$. with $m_1 \leq m_2$. Note that for all the possible values of θ the lighter neutrino mass eigenstate χ_1 consists of more left-handed neutrino than the right-handed one, which means that χ_1 always has sizable coupling to the SM Z^0 boson. Therefore the LEP-II bound on the mass of stable neutrino is always valid for χ_1, which is insensitive to the mixing angle.

As the fermions in the first three generations and the fourth generation have different $U(1)_F$ charges, the fourth generation fermions cannot mix with the ones in the first three generations through Yukawa interactions. After the spontaneous breaking down of $U(1)_F$, there exists a residual Z_2 symmetry for the fourth generation fermions which protect the fourth neutrino χ_1 to be a stable particle if it is lighter than the fourth generation charged lepton e_4, which makes it a possible dark matter candidate.

The thermal relic density of χ_1 is related to its annihilation cross section at freeze out. We numerically calculate the cross sections for $\chi_1\chi_1$ annihilation into all the relevant final states using CalHEP 2.4 [22]. In Fig. 1, we show the quantity $r_\Omega \equiv \Omega_{\chi_1}/\Omega_{DM}$ the ratio of the relic density of χ_1 to the observed total DM relic density $\Omega_{DM}h^2 = 0.110 \pm 0.006$ [23] as function of the mass of χ_1 for different values of the mixing angle θ. The results show a significant dependence on the mixing angle θ. For smaller mixing angle θ the couplings between χ_1 and gauge bosons $W^\pm, Z$ are stronger, resulting in a smaller relic density. The results also clearly show that due to the large annihilation cross section, χ_1 cannot make up the whole DM in the Universe. χ_1 can contribute to $\sim 20 - 40\%$ of the total DM relic density when its mass is around 80 GeV. But for $m_1 \gtrsim m_t$, it can contribute only a few percent or less to the whole DM.

However, since χ_1 has strong couplings to h^0 and Z^0, even in the case that the number density of χ_1 is very low in the DM halo, it is still possible that it can be detected by its elastic scattering off nucleus in direct detection experiments. Given the difficulties in detecting such a neutral and
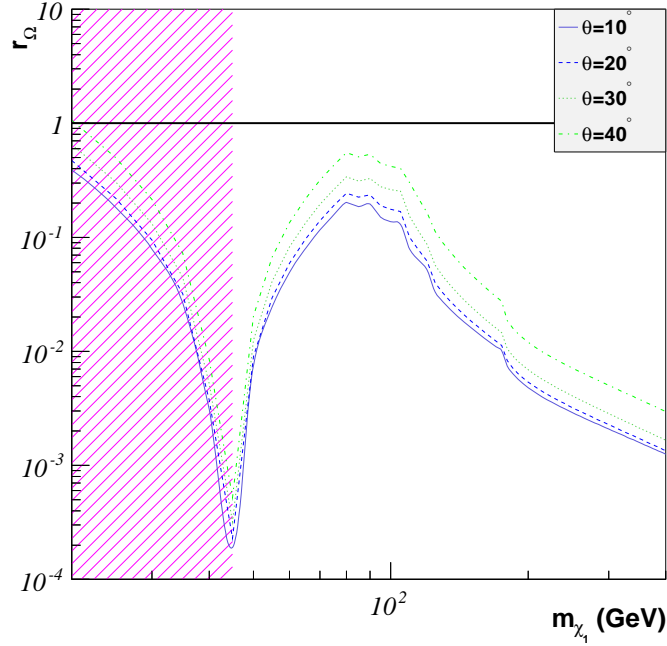
Figure 1: The rescaled χ_1 relic density r_Ω as function of the mass of χ_1. The shaded region is excluded by the LEP-II experiments.

stable particle at the LHC, there is a possibility that the stable fourth generation neutrino could be first seen at the DM direct detection experiments.

The generic formula for the differential event rate of DM-nucleus scattering per nucleus mass is given by

$$\frac{dN}{dE_R} = \frac{\rho_{DM}\sigma_N}{2m_{DM}\mu_N^2}F^2(E_R)\int_{v_{min}}^{v_{esc}}d^3v\frac{f(v)}{v},\tag{7}$$

where E_R is the recoil energy, σ_N is the scattering cross section corresponding to the zero momentum transfer, m_{DM} is the mass of the DM particle, $\mu_N = m_{DM}m_N/(m_{DM}+m_N)$ is the DM-nucleus reduced mass, $F(E_R)$ is the form factor, and $f(v)$ is the velocity distribution function of the halo DM. The local DM density ρ_{DM} is often set to be equal to $\rho_0 \simeq 0.3$ GeV/cm^3 which is the local DM density inferred from astrophysics based on a smooth halo profile. Since the neutrino DM can only contribute to a small fraction of the relic density of DM, it is likely that it also contributes to a small fraction of the halo DM density, namely, its local density ρ_1 is much smaller than ρ_0. We assume that ρ_1 is proportional to the relic density of χ_1 in the Universe, namely

$$r_\rho \equiv \frac{\rho_1}{\rho_0} \approx \frac{\Omega_{\chi_1}}{\Omega_{DM}},\tag{8}$$

or $r_\rho \approx r_\Omega$. Consequently, the expected event rates of the DM-nucleus elastic scattering will be scaled down by r_ρ. In order to directly compare the theoretical predictions with the reported experimental upper limits which are often obtained under the assumption that the local DM particle

density is ρ_0, we shall calculate the rescaled elastic scattering cross section

$$\tilde{\sigma} \equiv r_\rho \sigma \approx r_\Omega \sigma, \tag{9}$$

which corresponds to the event rate to be seen at the direct detection experiments. Note that $\tilde{\sigma}$ depends on the mass of χ_1 through the ratio r_ρ even when σ is mass-independent. The spin-independent DM-nucleon elastic scattering cross section in the limit of zero momentum transfer is given by [24]

$$\sigma_n^{SI} = \frac{4\mu_n^2}{\pi} \frac{[Zf_p + (A-Z)f_n]^2}{A^2}, \tag{10}$$

where Z and $A-Z$ are the number of protons and neutrons within the target nucleus, respectively. $\mu_n = m_1 m_n/(m_1 + m_n)$ is the DM-nucleon reduced mass. The couplings between DM and the proton (neutron) read

$$f_{p(n)} = \sum_{q=u,d,s} f_{Tq}^{p(n)} a_q \frac{m_{p(n)}}{m_q} + \frac{2}{27} f_{TG}^{p(n)} \sum_{q=c,b,t} a_q \frac{m_{p(n)}}{m_q}, \tag{11}$$

with $f_{Tq}^{p(n)}$ the DM coupling to light quarks and $f_{TG}^{p(n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{p(n)}$. The coefficient a_q in the model is given by

$$a_q = c_\theta^2 \frac{m_1 m_q}{v_H^2 m_h^2}. \tag{12}$$

The value of a_q is proportional to m_1, thus larger elastic scattering cross section is expected for heavier χ_1. The quark mass m_q in the expression of a_q cancels the one in the expression of $f_{p(n)}$. Thus there is no quark mass dependence in the calculations.

In Fig. 2 we give the predicted spin-independent effective cross sections $\tilde{\sigma}_n^{SI}$ for the fourth generation neutrino elastic scattering off nucleon as function of its mass for different values of the mixing angle θ. One sees that even after the inclusion of the rescaling factor r_ρ, the current Xenon100 data can still rule out a stable fourth generation neutrino in the mass range 55 GeV $\lesssim$ $m_1 \lesssim 175$ GeV which corresponds to $r_\Omega \lesssim 1\%$. Thus the stable fourth generation neutrino must be heavier than the top quark, and can only contribute to a small fraction of the total DM relic density On the other hand, for $m_{\chi_1} \gtrsim 200$ GeV, the cross section does not decrease with m_{χ_1} increasing, and is nearly a constant $\tilde{\sigma}_n^{SI} \approx 1.5 \times 10^{-44}$cm^2 in the range 200 GeV $\lesssim m_{\chi_1} \lesssim 400$ GeV. This is due to the enhanced Yukawa coupling between the fourth generation neutrino and the Higgs boson which is proportional to m_{χ_1}, as it is shown in the expression of a_q. Similar conclusions are expected for other models in which DM particles interact with SM particles through Higgs portal, for instance, the singlet scalar DM in extensions of left-right symmetry model [25–28]. One can see from the Fig. 2 that the result is not sensitive to the mixing angle θ either, which is due to the compensation of the similar dependencies on θ in the relic density. For instance, the cross sections for the $W^\pm W^\mp$ and $Z^0 Z^0$ channel of $\chi\chi$ annihilation are proportional to c_θ^4, which compensates the θ-dependence in the a_q for the elastic scattering processes.

The Majorana neutrino DM contributes also to spin-dependent elastic scattering cross section through axial-vector interaction induced by the exchange of the Z^0 boson. At zero momentum
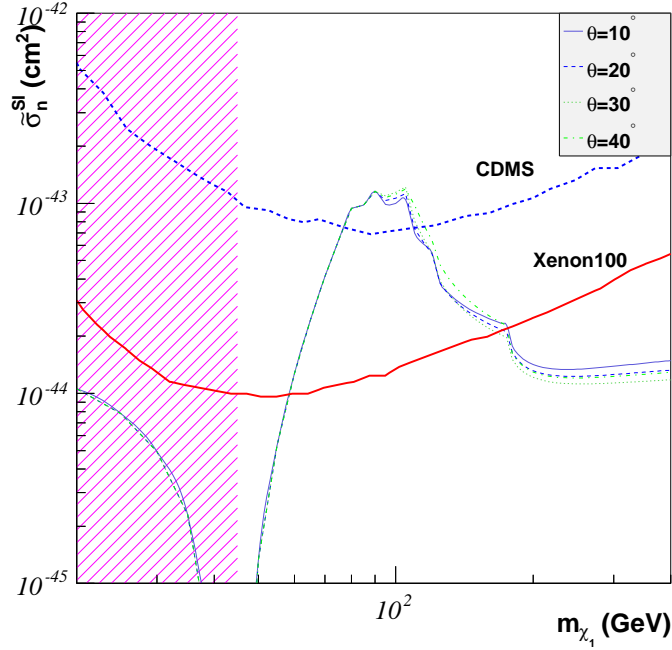
Figure 2: Effective spin-independent cross section $\tilde{\sigma}_n^{SI}$ which is σ_n^{SI} rescaled by $r_\rho \approx r_\Omega$ for χ_1 elastically scattering off nucleon as function of the mass of χ_1. Four curves correspond to the mixing angle $\theta = 10°$(solid), $20°$(dashed), $30°$(dotted) and $40°$(dot-dashed) respectively. The current upper limits from CDMS [29] and Xenon100 [30] experiments are also shown.

transfer, the spin-dependent cross section has the following form [24]

$$\sigma_N^{SD} = \frac{32}{\pi} G_F^2 \mu_n^2 \frac{J+1}{J} \left(a_p \langle S_p \rangle + a_n \langle S_n \rangle \right)^2,$$ (13)

where J is the spin of the nucleus, $a_{p(n)}$ is the DM effective coupling to proton (neutron) and $\langle S_{p(n)} \rangle$ the expectation value of the spin content of the nucleon within the nucleus. G_F is the Fermi constant. The coupling $a_{p(n)}$ can be written as

$$a_{p(n)} = \sum_{u,d,s} \frac{d_q}{\sqrt{2}G_F} \Delta_q^{p(n)},$$ (14)

where d_q is the DM coupling to quark and $\Delta_q^{p(n)}$ is the fraction of the proton (neutron) spin carried by a given quark q. The coefficients d_q in this model are given by

$$d_u = -d_d = -d_s = \frac{G_F}{\sqrt{2}}.$$ (15)

For the axial-vector interactions, the coupling strengths do not depend on the electromagnetic charges of the quarks.

In Fig. 3 we show the predicted effective spin-dependent DM-neutron cross section $\tilde{\sigma}_n^{SD}$ as function of the neutrino mass for different mixing angles, together with various experimental upper
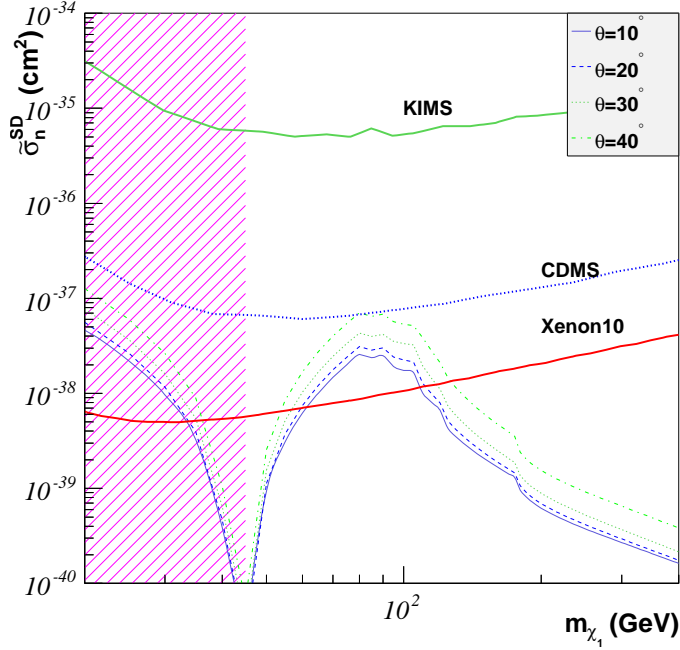
Figure 3: Effective spin-dependent cross section $\tilde{\sigma}_n^{SD}$ which is σ_n^{SD} rescaled by $r_\rho \approx r_\Omega$ for χ_1 elastically scattering off neutron as function of the mass of χ_1. Four curves correspond to the mixing angle $\theta = 10°$(solid), $20°$(dashed), $30°$(dotted) and $40°$(dot-dashed) respectively. The current upper limits from various experiments such as KIMS [34], CDMS [35] and Xenon10 [14] are also shown.

limits. Since σ_n^{SD} is independent of m_{χ_1}, the dependency of $\tilde{\sigma}_{p(n)}^{SD}$ on the neutrino mass comes from the dependency of r_ρ on m_{χ_1}, which can be seen by comparing Fig. 3 with Fig. 1. The Xenon10 data is able to exclude the neutrino DM in the mass range $60\,\mathrm{GeV} \lesssim m_{\chi_1} \lesssim 120\,\mathrm{GeV}$, which is not as strong as that from the Xenon100 data on spin-independent elastic scattering cross section. For a heavy neutrino DM with mass in the range $200\,\mathrm{GeV} \lesssim m_{\chi_1} \lesssim 400\,\mathrm{GeV}$ the predicted spin-dependent cross section is between $10^{-40}\,\mathrm{cm}^2$ and $10^{-39}\,\mathrm{cm}^2$.

In Fig. 4 we give the predicted spin-dependent DM-proton cross section $\tilde{\sigma}_p^{SD}$. The cross sections for Majorana neutrino DM scattering off proton and neutron are quite similar, which is due to the fact that the relative opposite signs in Δ_u and Δ_d are compensated by the opposite signs in d_u and d_n. So far the most stringent limit on the DM-proton spin-dependent cross section is reported by the SIMPLE experiment [31]. The SIMPLE result is able to exclude the mass range $50\,\mathrm{GeV} \lesssim m_{\chi_1} \lesssim 150\,\mathrm{GeV}$, which is compatible with the constraints from Xenon100. In Fig. 4, we also show the upper limits from indirect searches using up-going muons which are related to the annihilation of stable fourth generation neutrinos captured in the Sun. The limit from the Super-K experiment is obtained with the assumption that 80% of the DM annihilation products are from $b\bar{b}$, 10% from $c\bar{c}$ and 10% from $\tau\bar{\tau}$ respectively [32]. In the range $170\,\mathrm{GeV} \lesssim m_{\chi_1} \lesssim 400\,\mathrm{GeV}$, the limit from Super-K is $\sim 5 \times 10^{-39}\,\mathrm{cm}^2$. The IceCube sets a stronger limit $\tilde{\sigma}_p^{SD} \leq 2 \times 10^{-40}\,\mathrm{cm}^2$ for the DM mass at 250 GeV [33]. This limit is obtained with the assumption that the DM annihilation
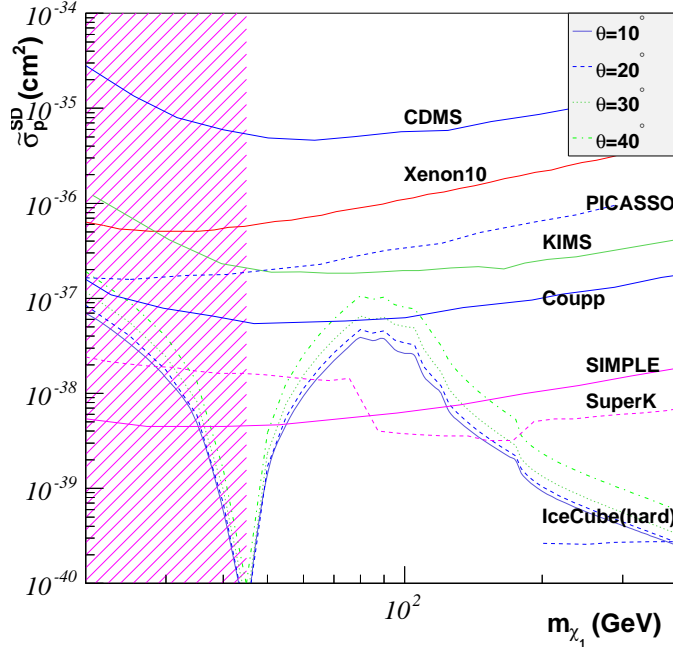
Figure 4: Effective spin-dependent cross section $\tilde{\sigma}_p^{SD}$ which is σ_p^{SD} rescaled by $r_\rho \approx r_\Omega$ for χ_1 elastically scattering off proton as function of the mass of χ_1. Four curves correspond to the mixing angle $\theta = 10°$(solid), $20°$(dashed), $30°$(dotted) and $40°$(dot-dashed) respectively. The current upper limits from various experiments such as KIMS [34], CDMS [35], Xenon10 [14], Coupp [36], Picasso [37], SIMPLE [38], SuperK [32], and IceCube [33] are also shown.

products are dominated by $W^\pm W^\mp$. If the annihilation products are dominated by $b\bar{b}$, the limit is much weaker, for instance $\tilde{\sigma}_p^{SD} \leq 5 \times 10^{-38}$ cm^2 for the DM mass at 500 GeV [33]. Note that in this model, the dominant final state is $Z^0 h^0$. The expected limit should be somewhere in between. Nevertheless, the IceCube has the potential to test these predictions.

In summary, we have investigated the properties of stable fourth generation Majorana neutrinos as dark matter particles. Although they contribute to a small fraction of the whole DM in the Universe, they can still be easily probed by the current direct detection experiments due to their relatively strong couplings to the SM particles. We have considered a fourth generation model with the stability of the fourth Majorana neutrino protected by an additional generation-dependent $U(1)$ gauge symmetry. We have shown that the current Xenon100 data constrain the mass of the stable Majorana neutrino to be greater than the mass of the top quark. For a stable Majorana neutrino heavier than the top quark, the effective spin-independent cross section for the elastic scattering off nucleon is found to be insensitive to the neutrino mass and is predicted to be around 10^{-44} cm^2, which can be reached by the direct DM search experiments in the near future. The predicted effective spin-dependent cross section for the heavy neutrino scattering off proton is in the range 10^{-40} cm$^2 \sim 10^{-39}$ cm^2, which can be tested by the ongoing DM indirect search experiments such as IceCube. This work is supported in part by the National Basic Research Program of China (973

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PoS(DSU 2012)028

11

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PoS(DSU 2012)028