If neutrinos are Dirac particles, then there will exist right–handed chirality states of neutrinos $\nu_R$ and left–handed chirality states of antineutrinos $\overline{\nu}_L$ which in the Standard Model Lagrangian do not have couplings to the weak interaction gauge bosons W and Z. The physical neutrino states are the helicity eigenstates $\nu_\pm$ and $\overline{\nu}_\pm$, and because helicity and chirality eigenstates do not coincide (only for ultra–relativistic neutrinos $\nu_- \simeq \nu_L$ and $\nu_+ \simeq \nu_R$), all the neutrino and antineutrino helicity states couple to the weak gauge bosons. Although the rate for W and Z mediated scattering between ultrarelativistic fermions involving one $\nu_+ (\overline{\nu}_-)$ is suppressed, depending on the process at hand, by a factor $(10^{-2} - 1)(m_\nu^2/2s)$, where $s$ is the centre–of–mass energy squared, as compared with the corresponding reaction involving $\nu_- (\overline{\nu}_+)$, the rate might be sufficiently large to thermalize these 'wrong–helicity' states in the early Universe.

If neutrinos with 'wrong' helicities should still be present with an energy density comparable to $\nu_-$ at the time when neutron–to–proton ratio froze out (at $T \simeq 0.7$ MeV), they would have speeded up the expansion rate of the Universe (which depends on the total energy density), with the consequence that the amount of produced primordial $^4$He would have exceeded the current observational limits. These imply that the effective number of extra relativistic two–component neutrino species $\delta N_\nu$ is much less than one, with $\delta N_\nu \lesssim 0.3$ perhaps the best limit [1]. Thus the production rate for $\nu_+$, and hence their mass [2], must be small enough so that they decouple already before the QCD phase transition which takes place at temperatures somewhere between 100 and 400 MeV. In that case they will not participate in the entropy transfer from quark–gluon plasma to particles in equilibrium, and consequently their number and energy densities will be diluted below levels that are acceptable for primordial nucleosynthesis. The production rate was first estimated by Fuller and Malaney [3], who argued that a Dirac neutrino with a lifetime exceeding the nucleosynthesis time scale ($t \sim 1$ s) should have a mass less than about 300 keV.

This argument applies to neutrinos with a mass much less than 1 MeV. A heavy neutrino with a mass in the MeV region would have a more pronounced effect on nucleosynthesis than a light neutrino, because during the synthesis of the
light elements the energy density of the 'right–helicity' states of a heavy neutrino would be comparable to or higher than that of a massless neutrino. This is because at that epoch these states have already decoupled \( T_{\text{dec}} \sim \text{few MeV} \). This has been shown \([4]\) to lead to an excluded region \(0.5 \text{ MeV} \lesssim m_{\nu_\tau} \lesssim 30 \text{ MeV} \) for the tau neutrino mass, provided \( \tau_{\nu_\tau} \gtrsim 10^3 \text{ s} \). (If \( 1 \text{ s} \lesssim \tau_{\nu_\tau} \lesssim 10^3 \text{ s} \), the upper bound is somewhat weakened.)

These considerations are relevant not only for the tau neutrino mass, the laboratory limit on which is \( m_{\nu_\tau} < 31 \text{ MeV} \) \([5]\), but possibly also for the muon neutrino. This is because the exprimental \( \nu_\mu \) mass limit has recently been revised upward \([6]\), with the current limit being about \( m_{\nu_\mu} < 500 \text{ keV} \).

The original Fuller and Malaney limit was based on an approximate estimate for the rate of \( \nu_+ \) production. In this Letter we shall present a careful re–evaluation of this limit by computing all the relevant cross sections and decay rates exactly. As we shall show by explicit calculation, the mass limit on \( \nu_\tau \) \( (\nu_\mu) \) becomes larger as much as by a factor of four (two) as compared with the Fuller and Malaney result.

Before the QCD phase transition but below, say, \( T \simeq 0.5 \text{ GeV} \), the fermions present in the Universe at significant number densities were the leptons and u, d, s and c quarks. All the \( 2 \rightarrow 2 \) scattering processes involving them, with no 'wrong–helicity' neutrinos in the initial state and with at least one \( \nu_+^\mu \) or \( \nu_+^\tau \) in the final state, are listed in Table 1. (There are altogether 47 separate reactions for each of the neutrinos which need to be taken into consideration). The first constraint is imposed on because each wrong helicity neutrino in the initial (final) state introduces an additional small factor \( m_1^2 / |p| \) \( (m_2^2 / |p'|) \) to the cross section. Here \( |p| \) and \( |p'| \) are the absolute values of the centre–of–mass momenta of the incoming and outgoing particles, respectively. Hence processes with more than one 'wrong–helicity' neutrino can be ignored as compared to processes with only one 'wrong–helicity' neutrino.

We have computed the cross sections for these scatterings in the limit of low momentum transfer \( |q^2| \ll M_Z^2, M_W^2, M_H^2 \). The contributions to the cross sections arising from the exchange of gauge bosons are given by

\[
\sigma_+^{(12 \rightarrow 34)} = \frac{G_F^2}{s} \frac{|p||p'|^3}{s} F_+^{(12 \rightarrow 34)} \left( \frac{m_1^2}{|p|^2}, \frac{m_2^2}{|p|^2}, \frac{m_3^2}{|p'|^2}, \frac{m_4^2}{|p'|^2} \right),
\]

(1)
where we have summed over all helicity and colour states of all charged leptons and quarks. The dimensionless functions $F^{(12\rightarrow34)}_\pm$ are listed in Table 1. In Eq. (1) the coefficient of $F^{(12\rightarrow34)}_\pm$ is of the order of an ordinary weak interaction cross section for fermion-fermion scattering, so that all suppression due to the 'wrong–helicity' neutrino in the final state is included in this function. For a massless neutrino $F^{(12\rightarrow34)}_\pm = 0$, and $F^{(12\rightarrow34)}_\pm$ increases monotonically as a function of its arguments. Some of the processes, such as $\nu_k^-\nu_k^+ \rightarrow \nu_k^+\nu_k^-$ and $\nu_k^+\nu_k^- \rightarrow \nu_k^+\nu_k^-$ ($k = \mu, \tau \neq j = e, \mu, \tau$) require two mass insertions and are therefore suppressed. Similar suppression holds in the relativistic limit also for the charged–current processes $l_k^+\nu_j^- \rightarrow \nu_k^+l_j^-$ and $l_k^+u_m \rightarrow \nu_k^+d_n$ ($u_m = u, c; d_n = d, s$), despite the fact that there is only one 'wrong–helicity' neutrino.

Having exact formulas one can estimate the cross sections for 'helicity–flip' scattering in the relativistic limit. It turns out that the popular approximation $\bar{\sigma}_\pm \approx G_F^2 E^2_\nu (m_\nu/2E_\nu)^2$, where the overbar indicates the averaging over the helicity and colour states of the initial charged leptons and quarks, overestimates the cross section for all the processes typically by more than an order of magnitude.

Neutral current scattering between fermions can also be mediated by the Higgs boson $H$. Although the Yukawa vertices provide direct 'helicity–flip' interactions for fermions, the fermion couplings to Higgs are weaker than gauge couplings roughly by a factor $m_f/M_W$. In the low energy limit additional suppression would arise from the propagators if the mass of the Higgs is larger than $M_W$ and $M_Z$. We have calculated the contributions from the Higgs boson exchange and from the interference between gauge boson(s) and the Higgs boson, and found that these are smaller at least by a factor $|p'|^2/M^2_H$ than gauge boson contributions to the cross section. In what follows we shall neglect the Higgs boson contribution.

Besides the scattering processes, $\nu_\mu$ and $\nu_\tau$ can be created also in three-body
decays. There exist six relevant decay channels producing $\nu_\mu$:

$$
\begin{align*}
\mu^- &\rightarrow \nu_\mu e^- \bar{\nu}_e, \\
\mu^- &\rightarrow \nu_\mu d \bar{u}, \\
c &\rightarrow s \mu^+ \nu_\mu^+, \\
c &\rightarrow d \mu^+ \nu_\mu^+, \\
\bar{s} &\rightarrow \bar{u} \mu^+ \nu_\mu^+, \\
\tau^+ &\rightarrow \nu_\tau^+ \mu^+ \nu_\mu^+,
\end{align*}
$$

(2)

and another six decay channels producing $\nu_\tau$:

$$
\begin{align*}
\tau^- &\rightarrow \nu_\tau^- e^- \bar{\nu}_e, \\
\tau^- &\rightarrow \nu_\tau^- \mu^- \bar{\nu}_\mu, \\
\tau^- &\rightarrow \nu_\tau^- d \bar{u}, \\
\tau^- &\rightarrow \nu_\tau^- \bar{d}, \\
\tau^- &\rightarrow \nu_\tau^- \bar{d} \\
\tau^- &\rightarrow \nu_\tau^- s \bar{c}, \\
\tau^- &\rightarrow \nu_\tau^- s \bar{u}.
\end{align*}
$$

(3)

The decay rate for particles, e.g. $\mu^- \rightarrow \nu_\mu^+ e^- \bar{\nu}_e$, is in the rest frame given by

$$
\Gamma_+^{(1\rightarrow 234)} = \frac{G_F^2}{12\pi^3} \int_0^{|p_4|_{\text{max}}} \frac{|p_4|^2}{\sqrt{|p_4|^2 + m_4^2}} \Delta \left(\frac{1}{2}\right) \frac{m_4^4}{m_4^{14}} \left\{ \Delta(m_{14}^2, m_2^2, m_3^2) + 2 \frac{m_4^4 (m_{14}^2 + m_2^2 + m_3^2)}{m_4^{14}} \left(m_1^2 \sqrt{|p_4|^2 + m_4^2} - m_1 |p_4| - m_2^2 \right) \right\} \frac{d|p_4|}{\sqrt{|p_4|^2 + m_4^2}} 
$$

(4a)

while the decay rate for antiparticles, e.g. $\tau^+ \rightarrow \nu_\tau^+ \mu^+ \nu_\mu^+$, is given by

$$
\Gamma_+^{(1\rightarrow 234)} = \frac{G_F^2}{6\pi^3} \int_0^{|p_4|_{\text{max}}} \frac{|p_4|^2}{\sqrt{|p_4|^2 + m_4^2}} \Delta \left(\frac{1}{2}\right) \frac{m_4^4}{m_4^{14}} \left\{ 2\Delta(m_{14}^2, m_2^2, m_3^2) + m_4^4 (m_2^2 + m_3^2) m_{14}^2 - 2(m_2^2 - m_3^2)^2 \right\},
$$

(4b)
with \( m_{14}^2 \equiv m_1^2 - 2m_1 \sqrt{|p_4|^2 + m_4^2} \) and \( \Delta(a,b,c) \equiv a^2 + b^2 + c^2 - 2ab - 2bc - 2ca \). Here the subscript 4 refers to \( \nu_4^\mu \) or \( \nu_4^\tau \) and the maximal momentum of the 'wrong–helicity' neutrino is \( |p_4|_{\text{max}} = \sqrt{|m_1^2 - (m_2 + m_3)^2 + m_4^2|^2 - 4m_1^2m_4^2/2m_1} \). The expressions for particle and antiparticle decay rates are different because they do not correspond to CP–conjugated processes: in both cases the final state involves the 'wrong–helicity' neutrino (but not anti–neutrino). In these rates we have summed over all helicity states of all particles except the \( \nu_+^{\mu(\tau)} \) under consideration.

If quarks are included in the decay process these decay rates must be multiplied by a factor \( 3V_{mn}^2 \) which accounts for the three colour states of quarks and for quark mixing.

The thermally averaged scattering rate reads [7]

\[
\Gamma_{\text{sc}}^+ = \frac{1}{n_{\nu_+^e}(T)} \sum_{(12\rightarrow34)} \int \frac{d^3p_1}{(2\pi)^3} \frac{d^3p_2}{(2\pi)^3} f(E_1/T) f(E_2/T) \sigma_{\text{sc}}^{(12\rightarrow34)} j(p_1,p_2), \tag{5}
\]

where \( n_{\nu_+^e} \) is the equilibrium number density of \( \nu_+^e \)’s, \( f(E_i/T) \) are the Fermi–Dirac distributions of the incoming particles, and \( j(p_1,p_2) = \sqrt{(p_1 \cdot p_2)^2 - m_1^2m_2^2/E_1E_2} \) is a flux–related factor. In (5) we have neglected the final state Pauli blocking, which is an about 10% effect [7].

Thermally averaged decay rate is simply given by

\[
\Gamma_{\text{d}}^+ = \frac{1}{n_{\nu_+^e}(T)} \sum_{(1\rightarrow234)} \Gamma_{\text{d}}^{(1\rightarrow234)} \int \frac{d^3p_1}{(2\pi)^3} f(E_1/T) \frac{m_1}{E_1}, \tag{6}
\]

where the factor \( m_1/E_1 \) arises from the Lorentz boost of the decay rate.

We have estimated the total thermally averaged \( \nu_+^e \) production rate \( \Gamma_+ = \Gamma_{\text{sc}}^+ + \Gamma_{\text{d}}^+ \) numerically, and the result is displayed in Fig. 1. The difference between the production rates for \( \nu_+^\mu \) and \( \nu_+^\tau \) is due to the differences in the phase spaces (initial and/or final) of the corresponding processes. While \( \mu \) is relativistic in the range of temperatures 100 – 400 MeV, \( \tau \) is nonrelativistic and therefore its number density, being Boltzmann–suppressed, is very sensitive to changes in temperature. This can be illustrated by the fact that while at \( T \approx 100 \text{ MeV} \) purely charged processes do not contribute to the \( \nu_+^\tau \) production, their contribution is about 40% at \( T \approx 250 \text{ MeV} \) and about 75% at \( T \approx 400 \text{ MeV} \). In contrast, for \( \nu_+^\mu \)
this figure rises from 55% \((T \approx 100 \text{ MeV})\) to 70% \((T \approx 400 \text{ MeV})\). One can see that elastic scattering and annihilation processes, dominated by \(\nu^k \bar{\nu}^k \rightarrow \nu^k \bar{\nu}^k\), are not so important as one would naively expect. Comparing with the Fuller and Malaney result, which is also shown in Fig. 1, we find that at \(T \approx 100 \text{ MeV}\) the actual rate for \(\nu^\tau_+ (\nu^\mu_+)\) production is 18 (7) times smaller.

We now require that the production of \(\nu^+\)'s ceases latest at the onset of the QCD phase transition, or that

\[
\Gamma_+(m_\nu, T_{\text{QCD}}) \leq H(T_{\text{QCD}}), \tag{7}
\]

where the Hubble parameter \(H\) is given by

\[
H = \sqrt{\frac{8\pi}{3M_{Pl}^2}} \rho(m_i, T) \equiv \sqrt{\frac{4\pi^3 g_{\text{eff}}(m_i, T)}{45}} \frac{T^2}{M_{Pl}}. \tag{8}
\]

Here \(\rho(m_i, T)\) is the total energy density, including also the particles that could be non-relativistic at the time of QCD phase transition, such as \(\tau, c, s\) and \(\mu\). Here we differ from the treatment of Fuller and Malaney, who did not account for the non-relativistic degrees of freedom in their estimate. Thus \(g_{\text{eff}}\) counts all the degrees of freedom, and we have tabulated it in Table 2, assuming that the non-relativistic species stay in equilibrium with the appropriate Boltzmann suppressed equilibrium densities. We have used the current quark masses \(m_u \approx 5.6 \pm 1.1 \text{ MeV}, m_d \approx 9.9 \pm 1.1 \text{ MeV}, m_s \approx 199 \pm 33 \text{ MeV},\) and \(m_c \approx 1.35 \pm 0.05 \text{ GeV} \ [8].\) About the half of the difference between the minimal and maximal values of \(g_{\text{eff}}\) is due to the uncertainty in quark masses. Another half comes from the uncertainty as to whether both or only one of the neutrinos \(\nu^\mu_+\) and \(\nu^\tau_+\) is in equilibrium.

As the production rate for the 'wrong–helicity' neutrino increases together with its mass (at \(T \gg m_\nu\ \Gamma_+ \propto m_\nu^2\)), and for the temperatures and masses under consideration \(H\) does not depend on \(m_\nu\), there exists an upper limit on the neutrino mass at which the inequality (7) can be satisfied. These mass limits for \(\nu_\tau\) and \(\nu_\mu\) as functions of the QCD phase transition temperature are displayed in Fig. 2. The value of the upper limit of the neutrino mass is almost insensitive to the
variations of the quark masses within ranges given above, despite the fact that the number and energy densities of s and c quarks change considerably. This happens because the changes in the production rate and the expansion rate compensate for each other. From Fig. 2 one can readily see that if $T_{\text{QCD}} \simeq (100)200$ MeV, the mass limits are $m_{\nu_\tau} \lesssim (1180)740$ keV and $m_{\nu_\mu} \lesssim (720)480$ keV. According to ref. [4], however, such large masses are already in conflict with primordial nucleosynthesis. Thus we may conclude that the equilibration of 'wrong' Dirac neutrino helicity states does not yield any additional new limit. Note that because the dilution provided by the entropy production at the QCD phase transition is more than sufficient for nucleosynthesis, $\nu_\tau$'s could well decouple some time during the actual phase transition, lowering the decoupling temperature and increasing the cosmological neutrino mass limit. Assuming decoupling at the onset of the phase transition will therefore slightly underestimate the actual upper limit.
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\[ F^{(12\rightarrow34)}_+(x_1, x_2, y_3, y_4) \]

| Process | \[ F^{(12\rightarrow34)}_+(x_1, x_2, y_3, y_4) \] |
|---------|----------------------------------|
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( x_{\nu^k}^2 \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 1/3 x_{\nu^k} (\sqrt{1 + x_{\nu^k}} + 1)^2 \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 1/2 x_{\nu^k} x_{\nu^k} \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( (1/12) x_{\nu^k} (\sqrt{1 + x_{\nu^k}} + 1)^2 \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 1/6 x_{\nu^k} [2(c_{V}\nu^k + c_{A}^A_{\nu^k})^2 + 3(c_{V}^2 + 3c_{A}^2_{\nu^k})x_{\nu^k}] \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 1/2 x_{\nu^k} [2(c_{V}\nu^k + c_{A}^A_{\nu^k})^2 + 3(c_{V}^2 + 3c_{A}^2_{\nu^k})x_{\nu^k}] \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 1/2 x_{\nu^k} [2(c_{V}\nu^k + c_{A}^A_{\nu^k})^2 + 3(c_{V}^2 + 3c_{A}^2_{\nu^k})x_{\nu^k}] \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( (1/12) y_{\nu^k} (\sqrt{1 + x_{\nu^k}} + 1)^2 \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( (1/3) y_{\nu^k} (c_{V}^2 + c_{A}^A_{\nu^k})(2 + 3 x_{\nu^k}) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( (1/3) y_{\nu^k} (c_{V}^2 + c_{A}^A_{\nu^k})(2 + 3 x_{\nu^k}) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 2(\sqrt{1 + y_{\nu^k}} - 1)(\sqrt{1 + x_{\nu^k}} + 1)(\sqrt{1 + y_{\nu^k}} - 1) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 2(3\sqrt{1 + x_{\nu^k}} - 1)(\sqrt{1 + y_{\nu^k}} - 1) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 2(3\sqrt{1 + x_{\nu^k}} - 1)(\sqrt{1 + y_{\nu^k}} - 1) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 12V^2_{\mu \mu}(\sqrt{1 + x_{\nu^k}} - 1)(\sqrt{1 + y_{\nu^k}} - 1)(\sqrt{1 + x_{\nu^k}} - 1) \) |
| \( \nu^k_\mu \nu^k_\mu \rightarrow \nu^k_\mu \nu^k_\mu \) | \( 4V^2_{\mu \mu}(\sqrt{1 + x_{\nu^k}} - 1)(\sqrt{1 + y_{\nu^k}} - 1)(\sqrt{1 + x_{\nu^k}} - 1) \) |

**Table 1.** List of 2 \( \rightarrow 2 \) scattering cross sections for \( \nu^k_\mu \) production. \( F^{(12\rightarrow34)}_+ \) is defined by Eq. (1), and the notations are: \( k = \mu, \tau \neq j = e, \mu, \tau; \ u^m = u, c; \ d^m = d, s; \ x_f \equiv m_f^2/|p|^2 \) and \( y_f \equiv m_f^2/|p|^2; \ c_{V}^f \equiv c_{V}^f + 1, \ c_{A}^f \equiv c_{A}^f + 1, \ c_{V}^f \equiv T^3_f - 2Q^f \sin^2 \theta_W \) and \( c_{A}^f \equiv T^3_f \) are the vector and axial vector couplings; \( V_{mn} \) is an element of the C-K-M matrix. We have used the values \( \sin^2 \theta_W = 0.2325, \ V_{ud} = 0.9753, \ V_{us} = 0.221, \ V_{cd} = 0.221 \) and \( V_{cs} = 0.9743 \) [8]. The present experimental uncertainties in these values do not affect our upper limits on the neutrino masses.
| $T$(MeV) | $g_{eff}^{min}$ | $g_{eff}^{max}$ |
|---------|----------------|---------------|
| 100     | 59.4           | 62.8          |
| 150     | 61.6           | 64.3          |
| 200     | 62.8           | 65.3          |
| 250     | 64.0           | 66.5          |
| 300     | 65.3           | 67.8          |
| 350     | 66.7           | 69.1          |
| 400     | 67.9           | 70.3          |

**Table 2.** The effective number of degrees of freedom as a function of temperature.
Figure Captions

**Fig. 1.** The total thermally averaged production rate for $\nu_+^\mu$ and $\nu_-^\tau$. For comparison, we show also the Fuller and Malaney rate (dashed curve).

**Fig. 2.** Neutrino mass limits corresponding to $\Gamma_+ \leq H$. The allowed region is below the curves.
Cosmological neutrino mass limit revisited

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Abstract

We consider the equilibration of the ‘wrong–helicity’ Dirac neutrino states $\nu_+$ and $\bar{\nu}_-$ in the early Universe via weak interactions and calculate carefully the thermally averaged production rate, taking into account all the relevant scattering and decay processes. Requiring that the production rate is less than the Hubble parameter at the onset of QCD phase transition so that the nucleosynthesis predictions are not contradicted, we find for $T_{\text{QCD}} \simeq 200$ MeV the upper limits $m_{\nu_e} \lesssim 740$ keV and $m_{\nu_\mu} \lesssim 480$ keV.

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