A Warm-Plus-Hot Dark Matter Universe

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Abstract

We investigate a new hybrid-model universe containing two types of dark matter, one “warm” and the other “hot”. The hot component is an ordinary light neutrino with mass \( \sim 25h^2 \) eV while the warm component is a sterile neutrino with mass \( \sim 700h^2 \) eV. The two types of dark matter arise entirely within the neutrino sector and do not require separate physical origins. We calculate the linear transfer functions for a representative sample of warm-plus-hot models. The transfer functions, and results from several observational tests of structure formation, are compared with those for the cold-plus-hot models that have been studied extensively in the literature. On
the basis of these tests, we conclude that warm-plus-hot dark matter is essentially indistinguishable from cold-plus-hot dark matter, and therefore provides a viable scenario for large scale structure. We demonstrate that a neutrino mass matrix can be constructed which provides the requisite dark matter constituents, while remaining consistent with all cosmological bounds.

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I. INTRODUCTION

Growing evidence against both the cold dark matter (CDM) and hot dark matter (HDM) cosmological models has prompted researchers to turn to alternative scenarios. Mixed dark matter (MDM) cosmologies [1,2], in which the universe contains an admixture of HDM, CDM, and ordinary matter, have captured most of the attention. By properly tuning the amount of HDM in the mix (20% – 30% in most scenarios) one can find models with a balance of small ($5h^{-1}$ Mpc) and large ($25–50h^{-1}$ Mpc) scale power consistent with current observations [2] ($h$ is the Hubble constant today in units of $100 \text{ km sec}^{-1}\text{Mpc}^{-1}$). Recent reports [3] from the Los Alamos National Laboratory suggesting that $\nu_e$ and/or $\nu_\mu$ have mass in the few eV range (just what the MDM models demand) have heightened interest in these models.

The essential property of HDM is that relativistic dark matter particles free-stream out of high-density regions, and therefore fluctuations are damped on scales smaller than the horizon size at the epoch when the HDM becomes non-relativistic. For ordinary neutrinos, the canonical HDM candidate, the size of the first objects to form is $\lambda_{FS} = 13\text{ Mpc} (\Omega_\nu h^2)^{-1}$, where $\Omega_\nu h^2 = m_\nu/93\text{eV}$ is the energy density in dark matter divided by the critical density, and $m_\nu$ is the neutrino mass. In MDM models, the cold component clumps on all scales. By shifting some of the mass density from hot to cold matter, one boosts small-scale power relative to large-scale power. This is the essence of the MDM models.

The greatest weakness of MDM models is that they require two types of dark matter with comparable mass densities. The hot component is usually taken to be an ordinary neutrino while the cold component is thought to be one of the standard CDM candidates (e.g., the lightest SUSY particle or the axion). A single dark matter candidate requires “new physics” beyond the Standard Model of strong and electroweak interactions. At first glance, a second type of dark matter would require adding another new sector to the particle physics theory, with the unexplained coincidence that the two types of dark matter give comparable contributions to the total mass density of the universe. (See, however, ref [1] for an attempt
to provide a particle physics connection between hot neutrinos and cold axions.)

In this paper, we explore the possibility that both the hot and cold dark matter components of a MDM universe are neutrinos. We propose that one of the active (weak interactions with matter) neutrinos is the hot component of the dark matter while one of the sterile neutrinos (no weak interactions) is a warm component. For the sterile neutrinos not to be hot, their mass to temperature ratio, $m_s/T_s$, must be greater than that of ordinary HDM. We therefore require that the sterile neutrinos decouple from the rest of the plasma prior to the electroweak (EW) phase transition. During the subsequent EW and QCD phase transitions, the sterile neutrinos cool relative to the coupled plasma (including the active neutrinos). $T_s$ is lowered and, assuming a fixed energy density, $m_s$ is higher. $m_s/T_s$ will be roughly 15 times higher than that of a standard ($\sim 25h^2$ eV) HDM neutrino. This type of particle is usually referred to as warm dark matter (WDM) and was first considered in the early 1980’s [4].

Large scale structure in a pure WDM universe with $m/T$ this large is much like large scale structure in a CDM-dominated universe. Even if the sterile neutrinos decouple between the QCD and EW phase transitions ($m_s/T_s$ roughly 8 times that of standard HDM neutrinos) they still behave much like cold dark matter [5].

Our main motivation is to show that two types of dark matter can arise simultaneously from the neutrino sector and provide a viable scenario for structure formation. We consider it a positive feature of the scenario that the two types of dark matter have masses which differ by only an order of magnitude or two. However, readers familiar with the standard lore of neutrino mass generation, will recognize our neutrino mass spectrum as unconventional. In the usual see-saw mechanism for generating small masses for active neutrinos, right-handed fields are typically very heavy due to the presence of large Majorana mass terms $M_R$ ($\sim 10^{3-19}$ GeV). The Dirac mass terms, which couple left and right-handed fields, are assumed to be comparable to the masses of the associated charged leptons or up-type quarks, i.e. $m_D \sim$ MeV-GeV. Assuming no left-handed Majorana mass terms, the physical eigenstates corresponding to the active neutrinos have masses $\sim (m_D^2/M_R)$, provided $m_D \ll M_R$. While the see-saw mechanism may be the most attractive scenario for explaining small (but non-
zero) masses for the active neutrinos, there is no \textit{a priori} reason for all three right-handed fields to be so heavy. The models constructed later are counterexamples where the physical state corresponding to one of the sterile neutrino types (\textit{i.e.} mostly right-handed fields) has mass $m_s \sim 700h^2$ eV.

\section*{II. LARGE SCALE STRUCTURE WITH TWO TYPES OF NEUTRINO}

Our model universe consists of four components: two types of massive neutrinos, two massless neutrino species that are treated as a single component, and a mixture of photons, baryons, and electrons that is treated as a single component ideal fluid. We assume primordial perturbations that are adiabatic, Gaussian and scale-free with a spectrum $P_p(k) = Bk^n$, where $k$ is the wavenumber of the perturbations measured in units of Mpc$^{-1}$. The COBE DMR experiment probes energy density perturbations on very large scales where there is little modification of the primordial spectrum. It finds $B = 8.2 \times 10^5h^{-4}$ Mpc$^4$ and estimates of $n$ that are consistent with $n = 1$ \cite{6,7}. We use this value of $B$ where normalization of the power spectrum is required and set $n = 1$. We also set $h = 0.5$ and ignore baryons. $h = 0.5$ is the value used in most studies of CDM and MDM \cite{2,8} though it may be in conflict with measurements of the Hubble constant \cite{9}. We emphasize that our main purpose here is to compare warm-plus-hot models with currently popular models such as MDM and not to carry out a detailed comparison with the observations.

Since the perturbations at early times are small, one can follow the initial stages of their evolution using linear theory. In fact, linear theory is adequate for studying structures on all scales significantly larger than $8h^{-1}$ Mpc. Quantitative tests on smaller scales probe non-linear structures, and therefore require N-body simulations.

The active neutrinos decouple from the photon-baryon plasma when they are still relativistic and when the photon temperature $T_\gamma \simeq 1$ MeV. Their background distribution function is therefore $f_0(p) = \left( e^{p/T_\nu} + 1 \right)^{-1}$ where $T_\nu = (4/11)^{1/3} T_\gamma$, $p = (E_\nu^2 - m_\nu^2)^{1/2}$, and $E_\nu$ is the neutrino energy. We assume that the background distribution function for the
sterile neutrino species also has a Fermi-Dirac shape; \( f_0(p) = \beta \left( e^{p/T_s} + 1 \right)^{-1} \). This distribution function is appropriate for a neutrino that decouples when it is relativistic provided \( \beta = 1 \) and \( T_s = (10.75/g_\ast)^{1/3}(4/11)^{1/3} T_{\gamma} \), where \( g_\ast \) is the effective number of relativistic degrees of freedom in the universe at the time of decoupling. The masses for the two neutrinos satisfy the relation

\[
m_\nu + m_s \left( \frac{T_s}{T_{\gamma}} \right)^3 = 93 \Omega_\nu h^2 \text{ eV}.
\]  

(1)

For a particle that decouples between the EW and QCD phase transitions, \( g_\ast \simeq 61 \), whereas a particle that decouples prior to the EW phase transition has \( g_\ast \simeq 107 \). We refer to models of these types as WPH1 and WPH2, respectively.

The above discussion assumes that the sterile neutrinos were coupled to the baryon-photon plasma at some early epoch. One possibility is that coupling occurs at very high energies where unknown physics (e.g. from Grand Unified Theories) operates. Alternatively, oscillations between active and sterile neutrinos may be responsible for bringing the sterile neutrinos into equilibrium [10,11]. In addition, Dodelson and Widrow [12] have shown that under certain assumptions, the momentum space distribution function for sterile neutrinos produced through oscillations has a Fermi-Dirac shape with \( T_s = (4/11)^{1/3} T_{\gamma} \), but \( \beta < 1 \). These models lead to exactly the same large scale structure phenomenology as the early-decoupled particle models [5] and so our discussion of structure formation applies to both.

It is convenient to define the linear transfer function \( T(k) \equiv (\delta \rho(k)/\rho)/(\delta \rho(k \rightarrow \infty)/\rho) \) which describes the evolution of the power spectrum through the linear regime. The evolved linear power spectrum is then \( P(k) = B k^n T(k)^2 \). The linear transfer functions for a variety of models are shown in Figure 1. As we can see, the WPH2 models are very close to the MDM models; the WPH1 models have a more distinct shape, though even in this case the differences are likely to be observationally indistinguishable. A more quantitative comparison of the models is made by calculating various integrals (with appropriate window functions) of the power spectrum. In particular, we calculate \( \sigma_8 \), EP (the excess power on 25 h\(^{-1}\)Mpc),
and $\sigma_{0.5}$. Here, $\sigma_L$ is the mass excess on $Lh^{-1}$ Mpc defined by the integral

$$
\sigma_L = \left( \int \frac{k^2 dk}{2\pi^2} P(k) W^2(kL) \right)^{1/2}
$$

(2)

where $W(x) = 3(\sin x - x \cos x) / x^3$ is the top hat window function. Fluctuations in the mass density are related to fluctuations in optically-selected galaxies through biasing, where $b_c\sigma_L$ gives the fluctuation in these galaxies on the scale $Lh^{-1}$ Mpc. $b_c$, the optical biasing parameter, can depend on scale, though in the simplest models it is assumed to be constant. Davis and Peebles [13] find that $b_c\sigma_8 \simeq 1$ and therefore $1/\sigma_8$ is a measure of the optical bias.

As mentioned above, CDM falls short because it has too little power on large scales relative to small scales. Wright et al. [6] introduce the quantity $EP \equiv 3.4\sigma_{25}/\sigma_8$. This definition is such that $EP = 1$ for standard CDM, whereas consistency with the APM angular correlation function requires $EP = 1.30 \pm 0.15$. Finally, $\sigma_{0.5}$ gives a rough (and probably overly pessimistic) approximation for $1 + z_{gf}$, where $z_{gf}$ is the redshift for galaxy formation. (Alternative methods for estimating $z_{gf}$ can be found in ref. [14] and references therein.)

The results for $\sigma_8$, EP, and $\sigma_{0.5}$ for a variety of models are given in Table 1. For MDM, CDM, and HDM, we use analytic fitting functions for $T(k)$ taken from Holtzman [15]. We see that the WPH1, WPH2, and MDM models with a 30% hot component give similar results for EP and $\sigma_8$. It has been argued [2] that MDM 30% satisfies large scale structure tests on $8 - 25h^{-1}$ Mpc scales, and therefore these conclusions should hold for our models as well. We do note, however, the results for $\sigma_{0.5}$ in Table 1, and also discussion in the literature (e.g. ref [2]), suggest that any model with a significant hot component may have trouble if observations push the redshift of galaxy formation $z_{gf} > 5$. This situation can be improved by decreasing the amount of hot matter in the mix, though this also leads to a decrease in power on large ($25h^{-1}$ Mpc) scales. Since galaxy formation depends on nonlinear and nongravitational physics (e.g., hydrodynamics) the above conclusions should be used with caution.
<table>
<thead>
<tr>
<th>Model</th>
<th>$m_\nu$</th>
<th>$m_s$</th>
<th>$\sigma_s$</th>
<th>EP</th>
<th>$\sigma_{0.5}$</th>
</tr>
</thead>
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<td>0.87</td>
<td>1.46</td>
<td>1.0</td>
</tr>
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<td>100</td>
<td>1.00</td>
<td>1.18</td>
<td>2.1</td>
</tr>
<tr>
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<td>90</td>
<td>0.93</td>
<td>1.27</td>
<td>1.6</td>
</tr>
<tr>
<td>WPH2 20%</td>
<td>4.6</td>
<td>190</td>
<td>1.02</td>
<td>1.17</td>
<td>2.6</td>
</tr>
<tr>
<td>WPH2 30%</td>
<td>7.0</td>
<td>160</td>
<td>0.94</td>
<td>1.26</td>
<td>2.0</td>
</tr>
<tr>
<td>MDM 30%</td>
<td>7.0</td>
<td>$\gg 10^3$</td>
<td>0.96</td>
<td>1.30</td>
<td>2.6</td>
</tr>
<tr>
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<td>$\gg 10^3$</td>
<td>1.48</td>
<td>0.94</td>
<td>9.6</td>
</tr>
</tbody>
</table>

Table 1 Results for $\sigma_s$, EP, and $\sigma_{0.5}$ for a variety of models. The results are based on linear theory calculations. Percentages in column 1 for the hybrid models indicate fraction of dark matter in the mix.

III. NEUTRINO MASSES – MODELS AND CONSTRAINTS

At temperatures comparable to $T_{QCD}$, sterile neutrinos interact with ordinary matter by mixing with active neutrinos. From detailed studies of neutrino oscillations in the early universe [10,11] the conditions that the sterile neutrinos are out of equilibrium below $T_{QCD}$ (and therefore do not spoil the predictions of Standard Big Bang Nucleosynthesis [16]) is

$$\delta m^2 < 3.6 \times 10^{-4}(\sin^2 2\theta)\ eV^2 \quad \text{for } \sin \theta_o \gtrsim 10^{-3} ,$$

where $\delta m^2 = m^2_\nu - m^2_s$ is the mass squared difference between the states, and $\theta_o$ is the vacuum mixing angle between them. This is provided that the energy of the states $E \gg \sqrt{\delta m^2}$, so that they do not decohere. However, for WPH1 it does not matter if oscillations bring the sterile neutrinos into equilibrium above $T_{QCD}$. As long as the sterile neutrinos decouple prior to the QCD phase transition, entropy transferred from the quark-gluon plasma into the interacting gas will dilute the sterile neutrino number density relative to the active neutrino number density. For this reason there is no constraint for $\sin \theta_o \lesssim 10^{-3}$ (The exact limit
depends on the adopted value for $T_{\text{QCD}}$). For WPH2, the demand that the sterile neutrinos are out of equilibrium below the EW phase transition at $T_{\text{EW}} \sim 300$ GeV, is essentially that of Eq. (3) except that the region where no constraint applies becomes $\sin \theta_0 \lesssim 10^{-8}$. (However, these bounds are model dependent. In the singlet majoron model [17], for example, they are greatly weakened due to the contribution of the majoron background to the active-neutrino self-energy [18].)

A specific mass matrix which is useful for further discussion is

$$
\begin{pmatrix}
0 & 0 & 0 & 0 & m_{e\beta} & m_{e\gamma} \\
0 & 0 & 0 & 0 & m_{\mu\beta} & m_{\mu\gamma} \\
0 & 0 & 0 & 0 & m_{\tau\beta} & m_{\tau\gamma} \\
0 & 0 & 0 & 0 & \mathcal{M}' & 0 & 0 \\
m_{e\beta} & m_{\mu\beta} & m_{\tau\beta} & 0 & M_{\beta\beta} & M_{\beta\gamma} \\
m_{e\gamma} & m_{\mu\gamma} & m_{\tau\gamma} & 0 & M_{\beta\gamma} & M_{\gamma\gamma}
\end{pmatrix},
$$

(4)

where we label our weak eigenstates as

$$
\nu_w = \left( \nu_{e(L)} \nu_{\mu(L)} \nu_{\tau(L)} \nu_{e(R)} \nu_{\mu(R)} \nu_{\tau(R)} \right).
$$

(5)

There is no reason to identify $\alpha$, $\beta$ and $\gamma$, with $e$, $\mu$ and $\tau$. $M_{\beta\gamma}$ is kept non-zero only to simplify later calculations. In (4) the matrix elements $m_{ij}$ and $M_{ij}$ correspond to Dirac and Majorana mass terms, respectively. As in most models, the upper left $3 \times 3$ submatrix of (4) is taken to be zero to be consistent with measurements of the width of the $Z_0$ (see however [19]). The other vanishing terms of (4) are set to zero for clarity, and in general they need not be identically zero for our scenario to remain viable (see later discussion).

Diagonalization of (4) leads to six mass eigenstates, which in the limit $m_{ij} \sim m \ll M_{ij} \sim M$ have mass eigenvalues

$$
m_1 = 0; \ m_2, m_3 = \mathcal{O}\left(\frac{m^2}{M}\right); \ m_4 = M'; \ m_5, m_6 \simeq M.
$$

(6)

For judicious choices of the entries in (4), the physical state corresponding to either $m_2$ or $m_3$ will correspond to the (active) HDM component and that corresponding to $m_4$ will be
the (sterile) WDM component. Those corresponding to \( m_5 \) and \( m_6 \) can be made sufficiently heavy with decay rates (into some combination of light neutrinos and possibly light scalars) short compared to the age of the Universe. The mass eigenstates corresponding to \( m_2 \) and \( m_3 \) will be a mixture of the \( \nu_{\alpha(L)} \), \( \nu_{\mu(L)} \) and \( \nu_{\tau(L)} \) states but with a small admixture \( \mathcal{O}(m/M) \) of the sterile \( \nu_{\beta(R)} \) and \( \nu_{\alpha(R)} \). If \( m/M \) were large enough, we see from (4) that the constraint \( m^2/M < 5 \times 10^{-3} \text{ eV} \) would be imposed. Since \( m^2/M \) is the mass of our light active state, this would imply that our HDM candidate could not possess a cosmologically interesting mass. However, if we take \( M \sim 1 \text{ TeV} \), then \( E < \sqrt{(b m^2)} \). As such, no oscillations will occur and we are relieved of this mass limit. The mixing angle with the lighter sterile state \( \nu_4 \) is zero so there is no oscillations into this state. From this we see that a mass matrix

with \( m \sim 3 \text{ MeV} \), \( M' \sim 200 \text{eV} \) and \( M \sim 1 \text{ TeV} \), would satisfy the oscillation constraints and yield the desired neutrino mass spectrum.

The mass matrix (4) has been chosen to elucidate our discussion. The first key feature of our scenario is the introduction of a new mass scale \( M' \ll M \). \( M \) may be the vacuum expectation value of one Higgs field and \( M' \) that of another. Alternatively, there may be only one new Higgs with different Yukawa couplings to the right-handed neutrino states. The smallness of \( M'/M \) is the main “unnatural” aspect of our scenario. However, to put it into perspective, the above-quoted value for \( M'/M \simeq 10^{-9} \) is not enormously smaller than some of Yukawa couplings in the Standard Model, such as \( g_Y (e) = m_e / M_{weak} \simeq 10^{-6} \), and is much larger than the ratio of Higgs vacuum expectation values postulated in Grand Unified Theories, \( M_W / M_{GUT} \simeq 10^{-14} \).

A second feature of our scenario is the non-participation (or at least reduced participation) of the light sterile state in the usual mixings of the left and right-handed neutrinos and in the consequent seesaw mechanism. In (4) this is accomplished by the introduction of matrix terms identically equal to zero. These zeros may be enforced by the imposition of global U(1) symmetries. For example, in the singlet majoron model [17] the Dirac mass terms arise from the coupling of the neutrinos to the the usual Standard Model Higgs doublet, \( H \). In addition, one introduces a complex scalar (Higgs) field \( \Phi \) which is a singlet under
SU(2)xU(1) gauge transformations, but carries non-zero Lepton-number. \( \Phi \) gets a vacuum expectation value generally taken to be \( \langle \Phi \rangle \approx \text{GeV} - 10 \text{ TeV} \). The right-handed Majorana masses come from their coupling to \( \langle \Phi \rangle \). If one extends this model by adding two new scalars \( \Phi_1 \) and \( \Phi_2 \) (or by adding another sterile neutrino), one can easily enforce the explicit zeros of (4) without generating any others – for example using the Lepton number assignment:

\[
\begin{array}{ccccccccc}
\nu_e(L) & \nu_\mu(L) & \nu_\tau(L) & \nu_e(R) & \nu_\mu(R) & \nu_\tau(R) & \Phi_1 & \Phi_2 & H \\
1 & 1 & 1 & 2 & 1 & 1 & -2 & -4 & 0
\end{array}
\]

In general the zeros of (4) need not be identically zero. That they be small relative to the other scales in the matrix would, in most cases, suffice. (The exact value required depends on the details of the model.) If indeed they are non-zero, then neutrino oscillations could be the mechanism that brings the sterile neutrinos into equilibrium at early times. There are two constraints we must impose on the non-zero mass terms. First, the admixture of active states in the heavier sterile mass eigenstate must be small enough that oscillations will not keep the sterile state in equilibrium below either \( T_{\text{EW}} \) (for WPH2) or \( T_{\text{QCD}} \) (for WPH1). (An explicit calculation in the context of the singlet majoron model shows that \( \nu_4 \) is decoupled prior to \( T_{\text{EW}} \).) Second, experimental limits on neutrino-less \( \beta\beta \)-decay require that \( 20 \)

\[
\langle m_{\nu_e} \rangle = |\sum U_{ei}^2 m_e| < 1.1 \text{ eV},
\]

where \( U_{ij} \) is the unitary matrix diagonalizing (4). This imposes limits on the values of \( m_{ij} \).

The discussion above provides an existence proof of our scenario. The mass matrix introduced and the constraints imposed on it are independent of the underlying particle physics model. However, there are other conditions we must satisfy, and to investigate them in any meaningful way requires the specification of a particular particle physics model. The model-dependent constraints are:

(i) The decay of the light sterile state \( \nu_s \) must be inhibited. The most general neutrino decay scheme can be written \( \nu_s \rightarrow \nu_a + X \) where \( \nu_a \) is some lighter neutrino (in our scenario one of the active states), and \( X \) is some set of bosons and/or pairs of lighter fermions.

(ii) Any new light degrees of freedom (e.g., the majoron) must be decoupled prior to \( T_{\text{QCD}}, \)
so as not to contribute significantly to the energy density of the universe at the epoch of nucleosynthesis. To ensure this, we assume that our new particle does not couple in any significant way with the active neutrino states or with ordinary matter. If it couples strongly only with the sterile neutrinos, then it to will be diluted away by entropy transfer at the electroweak and QCD phase transitions. Again, in the context of the singlet majoron model, this proves to be the case for the values of $m$, $M'$ and $M$ quoted.

(iii) The heavy degrees of freedom (here $\nu_{5,6}$) must annihilate or decay away before they can dominate the energy density of the universe. Again, in the singlet majoron model, the heavy states decay away on a sufficiently short time scale.

(iv) new $\beta\beta$-decay modes must be adequately suppressed. In the singlet majoron model, $g_{el}$ the effective coupling of the Majoron to $\nu_e$ must be $\leq 7 \times 10^{-5}$ [21]. This is satisfied by the model presented here.

The above constraints are clearly model dependent in that they rely on the details of the theory which ascribes mass to the neutrino sector. Any detailed investigation of the validity of our scenario therefore requires knowledge of the underlying particle physics. We have sometimes highlighted the possibilities for our scenario in the context of the singlet majoron model. Many other models possessing sterile neutrino fields, including more complicated majoron models, can also be constructed.

IV. CONCLUSION

We have investigated the possibility that the universe possesses a mixture of warm and hot dark matter. The hot component is identified with a known light neutrino with mass $\sim 25h^2$ eV; and the warm component with a sterile neutrino with mass $\sim 700h^2$ eV. These mass ranges satisfy the normal cosmological mass limits for stable neutrinos, provided the sterile neutrinos decouple at an early epoch. From calculations of the linear transfer functions for such a hybrid model, we show that mixed warm-plus-hot dark matter is a viable cosmological scenario, and in most respects provides a better fit to observational data than either the
standard HDM or CDM models. It is also interesting to note that the $700h^2$ eV sterile neutrino can evade the usual phase space constraints \cite{22}.

The introduction of our warm dark matter particle has come at a price, namely, the neglect of theoretical prejudice that all right-handed neutrino fields be very massive ($m > \mathrm{GeV}$). We emphasize, however, that \textit{a priori} the mass terms of matrix (4) are unknown, and there is no fundamental reason that sterile neutrinos cannot possess masses in the $700h^2$ eV range. In addition, we have discussed how such particles satisfy all standard cosmological bounds given certain restrictions on their oscillation properties. If theoretical prejudice is relaxed just a little, we see that a solution to the problems of the formation of large scale structure could reside purely in the neutrino sector.

We find it suggestive that one can accommodate a hybrid dark matter universe entirely within one sector avoiding the need for separate physical origins of the dark matter constituents. Given these arguments, and given the recent tentative experimental evidence for a hot neutrino component, we believe this model deserves further consideration.

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**Figure Caption**

Figure 1: Linear transfer functions for various models. The transfer functions for CDM, MDM, and HDM are taken from [15]. The hybrid models all assume 30% hot dark matter and 70% cold dark matter.