Late Time Neutrino Masses, the LSND Experiment, and the Cosmic Microwave Background

Z. Chacko,[1] Lawrence J. Hall,[1] Steven J. Oliver,[1] and Maxim Perelstein[2]
[1]*Department of Physics, University of California, Berkeley, Berkeley, California 94720, USA and Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA*
[2]*Institute for High-Energy Phenomenology, Cornell University, Ithaca, New York 14850, USA*
(Received 2 June 2004; published 23 March 2005)

Models with low-scale breaking of global symmetries in the neutrino sector provide an alternative to the seesaw mechanism for understanding why neutrinos are light. Such models can easily incorporate light sterile neutrinos required by the Liquid Scintillator Neutrino Detector experiment. Furthermore, the constraints on the sterile neutrino properties from nucleosynthesis and large-scale structure can be removed due to the nonconventional cosmological evolution of neutrino masses and densities. We present explicit, fully realistic supersymmetric models, and discuss the characteristic signatures predicted in the angular distributions of the cosmic microwave background.

DOI: 10.1103/PhysRevLett.94.111801 PACS numbers: 14.60.Pq, 12.60.Jv, 98.70.Vc, 98.80.Cq

Introduction.—The Liquid Scintillator Neutrino Detector (LSND) experiment found evidence for the oscillations $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ and $\nu_\mu \rightarrow \nu_e$ with an oscillation probability of around 3×10^{-3} [1] and a $\Delta m^2 \geq 1$ eV^2. The statistical evidence for the antineutrino oscillations is much stronger than for the neutrino case, with some analyses finding a 5σ effect [2]. While other experiments restrict the regions of parameter space that could explain the LSND data, they do not exclude the LSND result [3].

The confirmation of solar and atmospheric neutrino oscillations has led to a "standard" framework for neutrino masses, with three light neutrinos resulting from the seesaw mechanism and the heaviest right-handed neutrino not far from the scale of gauge coupling unification [4]. The LSND result conflicts with this framework, and, if confirmed by the MiniBooNE experiment [5], would throw neutrino physics into a revolutionary phase. There are three major challenges to incorporating the LSND result. A fourth light neutrino is needed with mass in the eV range, and this neutrino must be sterile. Such states are anathema to the seesaw mechanism. Secondly, neutrino oscillations in the early Universe would ensure that this fourth neutrino is thermally populated during big bang nucleosynthesis (BBN), yielding light element abundances in disagreement with observations [6]. In fact, theories with four neutrino species give very poor fits to the combined data of the LSND experiment and oscillation experiments with null results, while fits in theories with five neutrinos are much better [7]. This would imply that $N_{\nu\mathrm{BBN}} \geq 5$, in gross disagreement with observations [8,9]. Finally, the combination of large-scale structure surveys and Wilkinson Microwave Anisotropy Probe (WMAP) data [10,11] has led to a limit on the sum of the neutrino masses of about 0.7 eV, which is significantly less than the best fit values for the LSND neutrino masses [9,11].

An alternative explanation for why the neutrinos are light has been explored recently: the scale of neutrino masses can be dictated by a low-scale f of breaking of global symmetries [12,13]. Cosmologically, neutrinos remain massless until the symmetry-breaking phase transition, hence the name "late time neutrino masses". In this Letter, we argue that this scenario can accommodate the sterile neutrinos required by the LSND experiment. Moreover, the cosmological evolution of neutrino masses and densities in this scenario is nonstandard and, as a result, the apparent contradiction between the parameters preferred by LSND and cosmology can be avoided. At the same time, the scenario predicts potentially observable signatures in the cosmic microwave background (CMB).

All of the above features are generic in models with low-scale breaking of neutrino global symmetries, and can be understood without reference to a specific model.

The seesaw framework relies on the neutrinos' gauge quantum numbers to explain their masses and leaves only the active neutrinos light. Global symmetries may involve sterile and active neutrinos and may forbid Dirac and Majorana mass terms. Neutrino masses appear as a result of spontaneous breaking of these symmetries, so in this scenario, it is quite natural to expect light sterile states.

If the symmetry-breaking phase transition occurs *after* the BBN epoch, both active and sterile neutrinos are massless before and during nucleosynthesis. In this case, the oscillations that typically lead to thermal abundances for the sterile states in the traditional scenario are absent. During BBN, the energy density of the sterile neutrinos (and scalars required to break the global symmetries) is determined by their temperature. As we show below, this temperature can be significantly lower than that of the rest of the cosmic fluid. This allows our models to evade the BBN constraints on ΔN_ν.

The limit from large-scale structures on the sum of the neutrino masses is also easily avoided. The breaking of the global symmetries gives rise to a set of Goldstone bosons, which are coupled to both active and sterile neutrinos. This coupling is sufficiently strong for the sterile neutrinos to disappear after they become nonrelativistic, for example, by decaying into an active neutrino and a Goldstone boson. As a result, the relic abundance of the sterile neutrinos is low, and they do not significantly contribute to dark matter despite their large mass.

Specific Models.—There is a wide variety of models: the neutrinos may be either Dirac or Majorana, the number of sterile neutrinos may vary, and different choices of global symmetries and their breaking patterns can be made. Let us present two simple supersymmetric models incorporating the LSND neutrinos. For concreteness we construct models with three mass eigenstates that are predominantly sterile. We take the global symmetry to be $U(1) \times U(1)$; a simple possibility that allows a heavy neutrino to decay to light neutrino and a Goldstone boson.

Our first theory has three right-handed neutrino superfields, n. There is no overall lepton number symmetry, leading to six physical Majorana neutrinos. Above the weak scale the theory is described by the superpotential

$$\begin{aligned} W^M &= W_{\text{NMSSM}} + W_\nu^M, \\ W_\nu^M &= \lambda_{ij} l_i n_j h \frac{\phi}{M} + \frac{\kappa}{3}\phi^3 + \tilde{\lambda}_{ij} n_i n_j s \frac{\tilde{\phi}}{M} + \frac{\tilde{\kappa}}{3}\tilde{\phi}^3, \end{aligned} \tag{1}$$

where W_{NMSSM} is the superpotential of the nonminimal supersymmetric standard model (NMSSM); λ, $\tilde{\lambda}$, κ, and $\tilde{\kappa}$ are coupling constants; and the flavor indices i and j run from $1 \to 3$. The superfields l, h are the lepton and Higgs doublets of the MSSM, s is the electroweak singlet field of the NMSSM, and ϕ, $\tilde{\phi}$ are extra electroweak singlet fields. The nonrenormalizable operators in (1) are generated by integrating out physics at scale M; phenomenological constraints discussed below imply $M \sim 10^9$ GeV. In theories without an s field, the third operator in W_ν^M would be absent, and we would expect three light Dirac neutrinos. (If $nn\tilde{\phi}$ were allowed, the sterile states would be much heavier than the active states.) However, for theories such as the NMSSM, where the s field acquires a vacuum expectation value (VEV) of order the electroweak scale, the Dirac and Majorana mass terms are of the same order of magnitude, vf/M, explaining why the LSND neutrinos are quite close in mass to the active neutrinos. W_ν^M is the most general superpotential in the neutrino sector up to dimension four under the following discrete symmetries: Z_3, under which all the fields except ϕ and $\tilde{\phi}$ have charge $2\pi/3$; Z_3', under which s, h, and $\bar{h}$ are uncharged, q, l, n and $\phi, \tilde{\phi}$ have charge $2\pi/3$, while u^c, d^c and e^c have charge $-2\pi/3$; and Z_3'', under which n and $\tilde{\phi}$ both have charge $2\pi/3$ while ϕ has charge $-2\pi/3$.

Below the weak scale the renormalizable effective Lagrangian for the neutrino sector of the theory is

$$\mathcal{L}_\nu^M = g_{ij}\nu_i n_j \phi + \tilde{g}_{ij} n_i n_j \tilde{\phi} + \text{H.c.} + V(\phi, \tilde{\phi}), \tag{2}$$

where $g = \langle h \rangle \lambda / M$, $\tilde{g} = \langle s \rangle \tilde{\lambda}/\tilde{M}$ and the scalar potential is $V = -\mu^2|\phi|^2 + \kappa^2|\phi|^4 - \tilde{\mu}^2|\tilde{\phi}|^2 + \tilde{\kappa}^2|\tilde{\phi}|^4$. [We have assumed that supersymmetry (SUSY) breaking effects generate *negative* soft mass^2 terms for $\phi, \tilde{\phi}$.] This theory has two accidental $U(1)$ global symmetries: one under which ϕ and ν are charged and another one under which $\tilde{\phi}$, ν, and n are charged. When ϕ and $\tilde{\phi}$ acquire VEVs, these symmetries are broken leading to two pseudo Goldstone bosons G and $\tilde{G}$, and giving the neutrinos a mass.

With only minor changes we can construct a theory where the six neutrinos are Dirac. There are now three singlet left-handed sterile neutrinos, ν_i^s, and a total of six right-handed neutrinos n_α, coupled via

$$\begin{aligned} W^D &= W_{\text{NMSSM}} + W_\nu^D, \\ W_\nu^D &= \lambda_{i\alpha} l_i n_\alpha h \frac{\phi}{M} + \frac{\kappa}{3}\phi^3 + \tilde{\lambda}_{i\alpha}\nu_i^s n_\alpha s \frac{\tilde{\phi}}{M} + \frac{\tilde{\kappa}}{3}\tilde{\phi}^3. \end{aligned} \tag{3}$$

The superpotential W^D is the most general up to dimension four that is invariant under $Z_3 \times Z_3' \times Z_3''$ (with ν^s, like n, having charges $2\pi/3$ under each Z_3) together with lepton number symmetry.

Below the weak scale the renormalizable effective Lagrangian for the neutrino sector of this theory is

$$\mathcal{L}_\nu^D = g_{i\alpha}\nu_i n_\alpha \phi + \tilde{g}_{i\alpha}\nu_i^s n_\alpha \tilde{\phi} + \text{H.c.} + V(\phi, \tilde{\phi}). \tag{4}$$

The theory has two approximate global symmetries: one under which ϕ and ν are charged, and another under which $\tilde{\phi}$ and ν^s are charged. Again, ϕ and $\tilde{\phi}$ VEVs break these symmetries leading to two pseudo Goldstone bosons and Dirac masses for neutrinos.

Equations (2) and (4) imply that g and $\tilde{g}$ are of order m_ν/f, where m_ν is of order the neutrino masses, and f is the scale at which the global symmetries are broken. It is important to note that the couplings of the Goldstones bosons to the neutrinos are not diagonal in the neutrino mass basis. Instead, denoting the mass eigenstates by primes, these couplings are $(g_{\alpha\beta}\nu'_\alpha n'_\beta G + \tilde{g}_{\alpha\beta}\nu'_\alpha n'_\beta \tilde{G})$ (Dirac) and $(g_{\alpha\beta}\nu'_\alpha \nu'_\beta G + \tilde{g}_{\alpha\beta}\nu'_\alpha \nu'_\beta \tilde{G})$ (Majorana).

These theories provide concrete examples of a very rich set of theories. A particularly simple theory is obtained by deleting the $\tilde{\phi}$ field and removing the Z_3'' symmetry so that ϕ can couple to both doublet and singlet neutrino mass operators. In this case there is a single flavor diagonal $U(1)$ symmetry and a single Goldstone boson having diagonal couplings to neutrinos in the mass basis.

Constraints.— Significant constraints on the parameter space of these theories follow from the requirement that the total energy density in radiation during BBN does not differ significantly from the standard model. This requires that the "hidden sector" fields (ϕ, $\tilde{\phi}$, n, and ν^s, as well as the fermionic partners of ϕ and $\tilde{\phi}$ which will turn out to be quite light) not be in thermal equilibrium with the "visible sector" fields ($\nu, \gamma, \ldots$) before and during the BBN. We require that the two sectors decouple at a certain temperature $T_0 > 1$ GeV, and do not recouple until the temperature of the visible sector drops below $T_W \sim 1$ MeV, the temperature at which the weak interactions decouple. If this is the case, the reheating of the visible sector by the decoupling of heavy particles ($\mu, \pi, \ldots$) and possibly by the QCD phase transition will not affect the hidden sector. Defining r as the ratio of temperatures of the two sectors at the time of BBN, the energy density in the hidden sector is suppressed by a factor of r^4 compared to the naive estimate, and $r \lesssim 0.3$ allows one to avoid the BBN constraint even for a very large hidden sector.

The reactions that could recouple the two sectors include a $1 \leftrightarrow 2$ process $\phi \leftrightarrow \nu n$, $2 \leftrightarrow 2$ processes such as $\nu\bar{\nu} \leftrightarrow n\bar{n}$ and $\nu n \leftrightarrow \phi\tilde{\phi}$, $2 \leftrightarrow 3$ processes such as $\nu n \leftrightarrow 3\phi$, etc. Requiring that all these processes be "frozen" ($\Gamma < H$) prior to the weak interactions decoupling results in the following constraints on the couplings:

$$\begin{aligned} g_{ij}, g_{i\alpha} \lesssim 10^{-5}, \qquad g_{ij}\kappa, g_{i\alpha}\kappa \lesssim 10^{-10}r^{-1}, \\ g_{ij}\tilde{g}_{ij}, g_{i\alpha}\tilde{g}_{i\alpha} \lesssim 10^{-10}r^{-3/2}. \end{aligned} \tag{5}$$

Note that the coupling $\tilde{\kappa}$ is unconstrained.

The upper bounds on the coupling g can be translated into a lower bound on f. To interpret the LSND result in the model with Majorana sterile neutrinos, the low-energy theory must possess a mass term of the form $m_{ij}\nu_i n_j$, with at least some elements of m as large as 0.1 eV. This implies that $g_{ij}f \sim 0.1$ eV, and for a generic flavor structure we obtain a bound $f \gtrsim 10$ keV. A similar bound can be obtained for the Dirac neutrino case.

To avoid producing sterile neutrinos by oscillations prior to weak interactions decoupling, we require that the mass terms mixing active and sterile neutrinos not be generated until the temperature of the visible sector drops below T_W. Scattering in the plasma generates "thermal" masses for the ϕ bosons, $m^2(\phi) \sim \kappa^2 n_\phi(T')/T'$, where T' is the temperature of the hidden sector. The symmetry-breaking phase transitions for ϕ occur when $m^2(\phi) \sim \mu^2$. Using $\mu = f\kappa$, we conclude that the temperature of the visible sector at the time of this phase transition is $\sim f/r$, implying that $f \lesssim r$ MeV is necessary for the success of BBN. This in turn imposes a *lower* limit on the couplings, $g_{ij} \gtrsim r^{-1}10^{-7}$.

Thus, BBN considerations lead to an allowed range of

$$10 \text{ keV} \lesssim f \lesssim r \text{ MeV}. \tag{6}$$

Considerations of the supernova dynamics may slightly raise the lower bound; however, these constraints are strongly model dependent [14]. Even though the allowed values of f are much lower than the weak scale, the theory naturally allows for symmetry breaking in this range. The only assumption necessary is that ϕ only feels supersymmetry breaking through its couplings to l and n. Then μ^2 is of order $g^2 m_{\rm SUSY}^2/16\pi^2$, where $m_{\rm SUSY}$ is a typical soft supersymmetry breaking mass. Since g is of order m_ν/f and f itself is of order μ/κ, by eliminating μ and g in favor of f and m_ν we are led to the expression

$$f \approx \sqrt{\frac{m_\nu m_{\rm SUSY}}{4\pi\kappa}}. \tag{7}$$

For reasonable values of the parameters $m_\nu \approx 0.1$ eV, $m_{\rm SUSY} \approx 100$ GeV, $\kappa \approx 10^{-4}$, this yields f of order 3 MeV: quite close to the desired range. Analogous limits apply to the second symmetry-breaking scale, $\tilde{f}$.

Interestingly, the large-scale structure limit on the sum of neutrino masses [9] is *automatically* avoided in the models discussed here. The lower bound on g_{ij} obtained above implies that the reactions $\nu\bar{\nu} \leftrightarrow n\bar{n}, \phi\bar{\phi}$ become unfrozen *before* the sterile neutrinos become nonrelativistic. These reactions thermalize the hidden sector fields with the active neutrinos. The density of thermal, sterile neutrinos of mass m_s at temperatures $T < m_s$ is suppressed by a Boltzmann factor $e^{-m_s/T}$; the excess neutrinos either decay or annihilate. As a result, the massive sterile neutrinos do not make a significant contribution to dark matter. It is only the sum of the masses of active, stable neutrinos and the Goldstone bosons that has to satisfy the constraints of Ref. [10].

Signals in the CMB.—The nonstandard evolution of neutrino masses and densities in our scenario leaves an imprint in the CMB inhomogeneities [15]. There are two distinct, potentially observable effects [12]. First, the total relativistic energy density at the time of last scatter is modified due to the decay and annihilation of the sterile neutrinos. Second, unlike in the standard cosmology, free streaming of the active neutrinos may be prevented by their interactions with the Goldstone bosons.

During BBN, the energy density in the hidden sector is suppressed. When the reactions $\nu\bar{\nu} \leftrightarrow n\bar{n}, \phi\bar{\phi}$, $\nu n \leftrightarrow \phi$ become unfrozen, the two sectors thermalize and the energy density of the hidden sector fields approach that of the active neutrinos. (The active neutrinos themselves are by this time decoupled from electrons and photons.) These reactions, however, do not change the total relativistic energy density. When the sterile neutrinos become nonrelativistic ($T \sim m_s \sim 1$ eV) and decay or annihilate, the total relativistic energy is increased: this occurs at constant entropy, resulting in an increase in temperature. Since this occurs before matter-radiation equality, this will result in a nonstandard value of the relativistic energy density observed by CMB measurements. In terms of the "effective" number of neutrinos $N_{\nu,\rm CMB}$ [12], our scenario predicts

$$N_{\nu,\rm CMB} = 3\left(1 + \frac{n_s + 2.75 n_h/g_\nu}{3 + n_G/g_\nu}\right)^{1/3}. \tag{8}$$

Here, g_ν equals 7/4 for Majorana neutrinos and 7/2 for the Dirac case; n_h is the number of the "Higgs" (massive) components of the scalar fields that are light enough to be relativistic when the reactions $\nu\nu \to \phi\phi$, $\nu n \to \phi\tilde{\phi}$ be-

TABLE I. Effective number of neutrino species at last scatter, $N_{\nu,\mathrm{CMB}}$, as determined by the relativistic energy density.

n_G	Dirac n_s			Majorana n_s		
	1	2	3	1	2	3
2	3.59	3.78	3.95	3.78	3.92	4.06
3	3.70	3.86	4.01	3.91	4.03	4.14
8	4.00	4.11	4.21	4.22	4.29	4.35

come unfrozen; n_s is the number of sterile neutrino species, and n_G is the number of Goldstone modes. [Eq. (8) includes the contribution from the superpartners of the symmetry-breaking scalar fields.] For example, in the explicit models presented above, $n_s = 3$ and $n_G = n_h = 2$. Some typical values for $N_{\nu,\mathrm{CMB}}$ are presented in Table I. For comparison, while the current sensitivity on $N_{\nu,\mathrm{CMB}}$ from the WMAP and other CMB analyses [16] is about ± 5, the sensitivity of the Planck experiment is expected to reach the ± 0.20 level, providing a test of our predictions.

Furthermore, at the time of last scatter the mean free paths of the light neutrinos and the Goldstone bosons are well below the Hubble scale due to the process $\nu_i \leftrightarrow \nu_j G$. The absence of free-streaming leads to a shift in the positions of the CMB peaks at large l [12,17]. This shift (relative to the standard model prediction) is given by

$$\Delta l_n = 23.3 - 13.1\left(\frac{g_\nu(3 - n_S)}{(3g_\nu + n_G)(1/N_{\nu,\mathrm{CMB}} + .23)}\right) \quad (9)$$

where n_S is the number of light neutrinos that scatter during the eV era. (It is possible that $n_S < 3$ if one of the neutrinos is massless or very nearly so, or if m_G is large enough to make the process $\nu_i \rightarrow \nu_j G$ kinematically forbidden for some flavors.) Equation (9) provides another experimentally testable prediction of our scenario.

In the theory with no $\tilde{\phi}$ and one Goldstone boson, neutrino decays are absent so scattering only occurs via $\nu\nu \leftrightarrow GG$ and $\nu G \leftrightarrow \nu G$. The number of neutrino species which scatter is very sensitive to f and to whether the neutrino spectrum is hierarchical, inverted or degenerate.

M. P. is grateful to Alexander Friedland for useful discussions related to this work. Z. C., L. H., and S. O. are supported in part by the U.S. Department of Energy under Contracts No. DE-AC03-76SF00098 and No. DE-FG03-91ER-40676, and in part by the NSF grant no. PHY-00-98840. M. P. is supported by the NSF grant no. PHY-0098631.

Note added.—While completing this work the authors became aware of [18], where the bound on $\Sigma m_{\nu,i}$ is avoided by having heavy neutrinos annihilate to light scalars, increasing $N_{\nu,\mathrm{BBN}}$ significantly above 3.

[1] LSND Collaboration, C. Athanassopoulos *et al.*, Phys. Rev. Lett. **77**, 3082 (1996); Phys. Rev. C **58**, 2489 (1998); LSND Collaboration, A. Aguilar *et al.*, Phys. Rev. D **64**, 112007 (2001).

[2] A. Strumia, Phys. Lett. B **539**, 91 (2002).

[3] Particle Data Group Collaboration, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002).

[4] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity* (North-Holland, Amsterdam, 1979); T. Yanagida, in *Proceedings of Workshop on Unified Theories and Baryon Number in the Universe* (KEK, Tsukuba, 1979); R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[5] MiniBooNE Collaboration, R. Tayloe, Nucl. Phys. B, Proc. Suppl. **118**, 157 (2003).

[6] This conclusion can be avoided, however, in the presence of a non-negligible lepton asymmetry. See, for example, A. D. Dolgov, hep-ph/0306154, and references therein.

[7] M. Sorel, J. M. Conrad, and M. Shaevitz, Phys. Rev. D **70**, 073004 (2004).

[8] K. N. Abazajian, Astropart. Phys. **19**, 303 (2003); P. Di Bari, Phys. Rev. D **67**, 127301 (2003); Phys. Rev. D **65**, 043509 (2002).

[9] A. Pierce and H. Murayama, Phys. Lett. B **581**, 218 (2004); M. Cirelli, G. Marandella, A. Strumia, and F. Vissani, Nucl. Phys. B **708**, 215 (2005).

[10] D. N. Spergel *et al.*, Astrophys. J. Suppl. Ser. **148**, 175 (2003); W. J. Percival *et al.*, Mon. Not. R. Astron. Soc. **327**, 1297 (2001).

[11] A more conservative choice of priors can relax the bound on $\sum m_\nu$ to 1.7 eV; see SDSS Collaboration, M. Tegmark *et al.*, Phys. Rev. D **69**, 103501 (2004).

[12] Z. Chacko, L. J. Hall, T. Okui, and S. J. Oliver, Phys. Rev. D **70**, 085008 (2004).

[13] Models with global symmetry breaking at the *weak* scale have been proposed in Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Lett. B **98**, 265 (1981); G. Gelmini and M. Roncadelli, Phys. Lett. B **99**, 411 (1981); H. Georgi, S. Glashow, and S. Nussinov, Nucl. Phys. B **193**, 297 (1981). Global symmetries have been exploited to ensure the lightness of a sterile neutrino in R. N. Mohapatra, Phys. Rev. D **64**, 091301 (2001); B. Brahmachari, S. Choubey, and R. N. Mohapatra, Phys. Lett. B **536**, 94 (2002).

[14] See, for example, M. Kachelriess, R. Tomas, and J. W. F. Valle, Phys. Rev. D **62**, 023004 (2000).

[15] The possibility of lepton flavor violation as an additional, non-CMB, signal is also discussed in L. J. Hall and S. J. Oliver, Nucl. Phys. B, Proc. Suppl. **137**, 269 (2004).

[16] P. Crotty, J. Lesgourgues, and S. Pastor, Phys. Rev. D **67**, 123005 (2003); S. Hannestad, J. Cosmol. Astropart. Phys. 05 (2003), 004; E. Pierpaoli, Mon. Not. R. Astron. Soc. **342**, L63 (2003); V. Barger *et al.*, Phys. Lett. B **566**, 8 (2003).

[17] S. Bashinsky and U. Seljak, Phys. Rev. D **69**, 083002 (2004).

[18] J. Beacom, N. Bell, and S. Dodelson, Phys. Rev. Lett. **93**, 121302 (2004).