Realistic Dirac Leptogenesis

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We present a model of leptogenesis that preserves lepton number. The model maintains the important feature of more traditional leptogenesis scenarios: The decaying particles that provide the *CP* violation necessary for baryogenesis also provide the explanation for the smallness of the neutrino Yukawa couplings. This model clearly demonstrates that, contrary to conventional wisdom, neutrinos need not be Majorana in nature in order to help explain the baryon asymmetry of the universe.

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Introduction.—One interface between particle physics and cosmology is the attempt to provide an explanation for the observed baryon asymmetry in the universe. Leptogenesis represents one of the most attractive possibilities for the generation of this asymmetry. The recent discovery of neutrino masses has further increased the credibility of this scenario. In its original incarnation [1], leptogenesis relies upon the decay of right-handed Majorana neutrinos to create lepton number, which is subsequently transformed into baryon number by the electroweak $B + L$ anomaly. This traditional scenario relies in an essential way on the breaking of lepton number by the Majorana right-handed neutrinos. The attractive feature of this model is that the right-handed neutrinos responsible for the generation of the lepton asymmetry are also responsible for the smallness of the observed neutrino masses through the seesaw [2] mechanism.

Since the original model of (Majorana) leptogenesis, there have been two important observations. First, the provoking observation has been made that it is not necessary to break lepton number to have a theory of leptogenesis, and that leptogenesis could be accomplished in a theory with Dirac neutrinos [3]. We will review this idea in the next section. A disadvantage of this idea, relative to the traditional models of leptogenesis, is that it possesses no relationship between the mechanisms responsible for the generation of the lepton asymmetry and the smallness of the neutrino masses. The second observation was that, in supersymmetric theories, it is possible to explain the smallness of the neutrino Yukawa couplings by relating their presence to supersymmetry breaking [4]. Combining these two ideas allows us to once again relate the generation of the lepton asymmetry to the smallness of neutrino masses.

This brings Dirac leptogenesis onto a footing equal to that of the traditional Majorana leptogenesis models. The only ingredient that this mechanism requires beyond the usual leptogenesis scenario is a $\mathrm{U}(1)_N$ symmetry, which forbids the bare Yukawa couplings between the left- and right-handed neutrinos.

Review of leptogenesis with Dirac neutrinos.—Reference [3] noted that, even in a theory that conserves lepton number, a *CP* violating decay of a heavy particle can result in a nonzero lepton number for left-handed particles, and an equal and opposite nonzero lepton number for right-handed particles. For most standard model species, Yukawa interactions between the left-handed and right-handed particles are sufficiently strong to cancel these two stores of lepton number rapidly. However, the interactions of a right-handed Dirac neutrino are exceedingly weak, and equilibrium between left-handed lepton number and right-handed lepton number will not be reached until temperatures fall well below the weak scale. By this time lepton number has already been converted to baryon number by sphalerons.

To see how this scenario works, imagine that a negative lepton number is stored in the left-handed neutrinos, while a positive lepton number of equal magnitude is stored in the right-handed neutrinos. Sphalerons act only on left-handed particles, violating $B + L$ while conserving $B - L$. This means part of the negative lepton number stored in left-handed neutrinos can be converted to a positive baryon number by the electroweak anomaly. The (now smaller in magnitude) negative lepton number stored in the left-handed neutrinos ultimately equilibrates with the positive lepton number stored in the right-handed neutrinos only after the temperature of universe drops below electronvolts. The processes responsible for equilibrating the right- and left-handed neutrinos conserve both B and L separately. The ultimate result is a universe with a total positive lepton number and a total positive baryon number.

Small Yukawa couplings.—The basic program in this Letter is to generate small Dirac Yukawa couplings by integrating out a heavy field following the methods of [5]. The smallness of the Yukawa couplings will be explained by the large ratio between the scale of supersymmetry breaking and the heavy masses. The key point is that the same heavy fields can be responsible for the generation of the *CP* asymmetry.

The heavy fields to be integrated out are three pairs of vectorlike leptons, ϕ and $\bar{\phi}$, one pair per generation of standard model particles. These fields transform as doublets under $SU(2)_L$ [6]. The decay of these heavy leptons will also provide the necessary *CP* violation for leptogenesis. In order to have *CP* violation in this sector, it is sufficient to have two generations.

We work in the context of the minimal supersymmetric standard model (MSSM) augmented by three generations of right-handed neutrinos. We forbid bare Yukawa couplings, LNH_u, through the use of a $U(1)_N$ symmetry [7], under which the N has charge $+1$, while all the fields of the MSSM are uncharged. We also add a gauge singlet, χ that breaks $U(1)_N$ when it acquires a vacuum expectation value (vev). The field content of the model, along with the charges under $SU(2)_L$, $U(1)_N$, $U(1)_N$, and $U(1)_Y$ is shown in Table I. $U(1)_N$ is the standard lepton number, which remains a symmetry in this model broken only by the $SU(2)_L$ anomaly. With these charge assignments, the most general renormalizable superpotential is

$$\mathcal{W} \ni \lambda N\phi H_u + hL\bar{\phi}\chi + M^{\phi}\phi\bar{\phi}, \tag{1}$$

where we have suppressed generation indices. Upon integrating out the heavy vector lepton pair, we get the following superpotential:

$$\mathcal{W}_{\rm eff} \ni \lambda h\frac{NH_uL\chi}{M^{\phi}}. \tag{2}$$

Next, we arrange for the χ field to take on a weak-scale vev. We can accomplish this, for example, through the use of an O'Raifeartaigh model of the type used for neutrino masses in [9]. This approach gives $\langle F_\chi\rangle \simeq m_{3/2}M_{\rm Planck} \neq 0$ and $\langle\chi\rangle = 0$ in the limit of global supersymmetry, but $\langle\chi\rangle \simeq 16\pi^2 m_{3/2}/\kappa^3 \neq 0$, where κ is a dimensionless coupling constant, after supergravity effects are taken into account. Because of the large $\langle F_\chi\rangle$, left-handed and right-handed sneutrinos equilibrate quickly above the weak scale. However, the asymmetry stored in the right-handed neutrino (fermion) remains intact. Interesting collider phenomenology could result from the large F_χ [4]. In any case, it is clear that the Dirac neutrino Yukawa couplings, y_ν, will be suppressed by the ratio of the weak scale to the heavy masses:

$$y_\nu \sim h\lambda\frac{\langle\chi\rangle}{M^{\phi}}. \tag{3}$$

Because $\langle\chi\rangle$ does not have to be exactly at the electroweak scale, it gives an additional freedom beyond the traditional Majorana leptogenesis. We note that a very similar superpotential was considered in [10][10], with the vev of the χ field replaced with a hard mass.

Lepton asymmetry.—It remains to check whether this scenario can generate a sufficient baryon asymmetry. *CP* violation will enter the theory through the decay of the ϕ and $\bar{\phi}$ particles. There are equal contributions from the decay of the scalar and fermionic components. For simplicity, we will concentrate on the decay of the scalars. The leading order contribution to the *CP* violation in ϕ decay comes from the interference between the tree-level diagrams and the absorptive part of the one-loop wave function renormalization diagrams [11–13]. In the case where there are two "generations" of $\phi - \bar{\phi}$ pairs, it is possible to rotate away all but one physical phase. We will consider this case in the following. Additional generations of the $\phi - \bar{\phi}$ pairs will just allow for the possibility of additional baryon number generation. In addition, the two generation case is a good approximation if the masses of the ϕ particles are reasonably well separated. We take the mass matrix, M^{ϕ}, to be diagonal with elements M_1 and M_2. The diagrams relevant to the calculation of the *CP* asymmetry are shown in Fig. 1. Note that, in this model, the one-loop vertex renormalization diagrams do not violate *CP* and are therefore irrelevant to the

TABLE I. The field content and quantum numbers of the model.

Field	$U(1)_L$	$U(1)_N$	$SU(2)_L$	$U(1)_Y$
N	-1	$+1$	**1**	0
L	$+1$	0	**2**	$-\frac{1}{2}$
H_u	0	0	**2**	$\frac{1}{2}$
ϕ	$+1$	-1	**2**	$-\frac{1}{2}$
$\bar{\phi}$	-1	$+1$	**2**	$\frac{1}{2}$
χ	0	-1	**1**	0

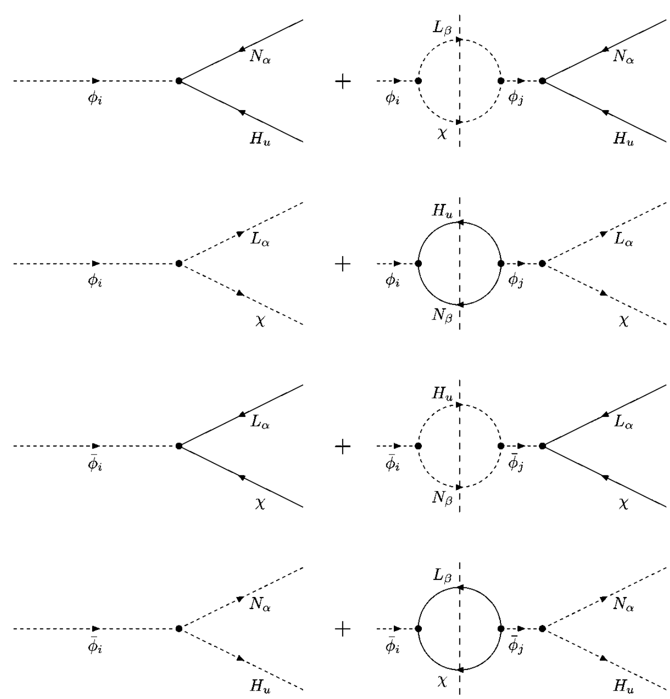

FIG. 1. Diagrams giving the leading contribution to the *CP* asymmetry in ϕ and $\bar{\phi}$ scalar decays. The absorptive part of the one-loop diagrams contributes to the *CP* asymmetry.

calculation of the lepton asymmetry. Restoring the generation indices to Eq. (1), we have

$$\mathcal{W} \ni \lambda_{i\alpha} N_\alpha \phi_i H_u + h_{\beta i} L_\beta \bar{\phi}_i \chi + M_a^\phi \phi_a \bar{\phi}_a, \quad (4)$$

Now we proceed with the calculation of the asymmetry. In the case where the magnitudes of the masses $|M_1|$ and $|M_2|$ are well separated, the asymmetry will be dominated by the decay of the lightest $\phi - \bar{\phi}$ pair (we take $|M_1| < |M_2|$) and can readily be calculated (following the methods of [13]). We now define the quantities $J \equiv \mathrm{Im}(h_{\beta 1}^* h_{\beta 2} \lambda_{1\alpha}^* \lambda_{2\alpha} M_1 M_2^*)$ and $\Delta M^2 \equiv |M_1|^2 - |M_2|^2$. In J, the α and β indices run over the generations of the L and N particles. For the decay asymmetries, we find

$$\epsilon_{\bar{N}} \equiv \frac{\Gamma(\phi_1 \to N^c H_u^c) - \Gamma(\phi_1^c \to N H_u)}{\Gamma(\phi_1)} = -\frac{J}{4\pi \Delta M^2 (|\lambda_{1\alpha}|^2 + |h_{\beta 1}|^2)} \equiv \varepsilon, \quad (5)$$

$$\epsilon_L \equiv \frac{\Gamma(\phi_1 \to L\chi) - \Gamma(\phi_1^c \to L^c \chi^c)}{\Gamma(\phi_1)} = -\varepsilon, \quad (6)$$

$$\epsilon_{\bar{L}} \equiv \frac{\Gamma(\bar{\phi}_1 \to L^c \chi^c) - \Gamma(\bar{\phi}_1^c \to L\chi)}{\Gamma(\bar{\phi}_1)} = \varepsilon, \quad (7)$$

$$\epsilon_N \equiv \frac{\Gamma(\bar{\phi}_1 \to N H_u) - \Gamma(\bar{\phi}_1^c \to N^c H_u^c)}{\Gamma(\bar{\phi}_1)} = -\varepsilon. \quad (8)$$

Note that $\Gamma(\phi) = \Gamma(\bar{\phi})$ due to supersymmetry, because chiral superfields ϕ and $\bar{\phi}$ form a massive supermultiplet. Here we have used the same names for fermion and scalar fields in the same multiplet, and the α and β indices labeling the generation of the final state particles are summed over. The above asymmetries in the decay amplitude give rise to a store of lepton number in the left-handed and right-handed (s)neutrinos. In the limit that the particles decay well out of equilibrium (the "drift and decay" limit), the asymmetry is given by [14]

$$N \equiv \frac{n_N}{s} \sim \frac{(\epsilon_N - \epsilon_{\bar{N}}) n_\gamma}{g_* n_\gamma} \sim \frac{-2\varepsilon}{g_*}, \quad (9)$$

$$L \equiv \frac{n_L}{s} \sim \frac{(\epsilon_L - \epsilon_{\bar{L}}) n_\gamma}{g_* n_\gamma} \sim \frac{-2\varepsilon}{g_*}. \quad (10)$$

However, this limit is not necessarily applicable, as the condition for out-of-equilibrium decay, $\Gamma(\phi_1)/2H(M_1) \lesssim 1$, is only marginally satisfied. Therefore, one should solve the full system of Boltzmann equations numerically, including $2 \to 2$ scattering, to accurately determine the lepton asymmetry. However, for an existence proof that this mechanism will work, we will not need to resort to these numerics: We simply note that for the specific choices of $\lambda = h^T$ and $\langle \chi \rangle$ equal to the electroweak vev, our asymmetry (and neutrino mass matrices) will reduce to that of the standard supersymmetric leptogenesis scenario with Majorana neutrinos. It has been shown (for recent reviews see [15]), that the generation of a sufficient lepton asymmetry is possible in this case, with the mass of heavy neutrinos at the 10^{10} GeV scale. Indeed, it is possible that more complicated textures for λ and h might lead to a more efficient generation of a lepton asymmetry while remaining consistent with low-energy data on neutrino oscillations.

Cosmological and astrophysical constraints.— Theories of supersymmetric leptogenesis have tension with the gravitino problem; the reheat temperature must be low enough to avoid cosmological difficulties associated with gravitino production. A typical constraint is $T_{RH} \lesssim 10^9$–10^{10} GeV for 1–2 TeV gravitino [16]. On the other hand, the reheat temperature must be high enough to produce the particles (in our case the ϕ and $\bar{\phi}$) that need to be heavy in order to decay out of equilibrium. However, as we have shown above, our scenario can reproduce a baryon asymmetry equal to that of the traditional leptogenesis scenario, which has been shown to be compatible with gravitino constraints [15]. There are a host of other ideas to help with this tension. For example, theories of anomaly mediation [17] have gravitino masses that are heavier than the usual case by a loop factor, of order 100 TeV. Furthermore, there has been recent work suggesting that it may be possible to significantly increase the mass of the gravitino in theories with weak-scale supersymmetry, thereby obviating the gravitino problem [18].

Yet another possibility involves using coherent oscillations of the scalar fields carrying lepton number [19,20]. In our case the $\phi = \bar{\phi}$ flat direction could be used, for example, with the O'Raifeartaigh model discussed earlier with $\kappa \sim 1$, $\langle \chi \rangle \sim 10$ TeV. We make the assumption that N and L remain pinned to the origin. If we stick to the simplifying ansatz $\lambda = h^T$, we can scale M^ϕ proportional to $\langle \chi \rangle$ so as to reproduce the observed neutrino masses with the same Yukawa couplings as the traditional case. This means that the CP asymmetry remains the same as well. Working within the model of [20] (replacing N with the $\phi = \bar{\phi}$ flat direction), in order to have the CP asymmetry large enough, we require $M_1^\phi \gtrsim 10^8$ GeV. This can well be consistent with the gravitino mass of $\sim$1 TeV. In addition, the possibility $\lambda \neq h^T$ gives even more freedom.

It would be interesting to study the gravitino problem with both $\langle \chi \rangle$ and $\langle F_\chi \rangle$ (and, hence, the gravitino mass) as free parameters, such as in models of gauge-mediated supersymmetry breaking. Smaller $\langle F_\chi \rangle$ gives a lighter gravitino, and the constraint on the reheat temperature is more severe [21]. However, smaller $\langle F_\chi \rangle$ allows smaller $\langle \chi \rangle$ while preventing the appearance of a negative eigenvalue in the sneutrino mass-squared matrix. This, in turn, would allow for lighter ϕ, which helps with the gravitino problem. Therefore, we expect Dirac leptogenesis to accommodate models with lower $\langle F_\chi \rangle$ more easily than traditional leptogenesis models.

There might be a worry that the right-handed neutrinos could potentially represent a dangerous number of additional light species at the time of big-bang nucleosynthesis (BBN). The constraint is $\Delta N_\nu \lesssim 0.3$ [22]. However, by the time of BBN, the contribution of right-handed neutrinos is suppressed by the entropy factor: $\Delta N_\nu = 3(T_{\nu_R}/T_{\rm bath})^4 = 3[g_*(1\ {\rm MeV})/g_*({\rm MSSM} + N)]^{4/3} = 0.02$ and is safe.

When the $U(1)_N$ symmetry is broken by the χ vev or F_χ vev, a Nambu-Goldstone boson will be produced. Generally, stringent astrophysical constraints on such particles (e.g., Majorons, familons) are derived from looking at supernovae. The usual constraints assume couplings between the SM fields and the Nambu-Goldstone bosons. In contrast, in this case the right-handed neutrino is the only light field charged under the $U(1)_N$. Consequently, the couplings of the Nambu-Goldstone bosons to the matter in the supernova will be exceedingly weak. Nambu-Goldstone boson production processes will be suppressed by factors of m_ν/T relative to the usual case. Since even the usual case (see, for example, [23]), can be made acceptable, there is clearly no problem here.

Conclusion.—We have presented a realistic model of supersymmetric leptogenesis using Dirac neutrinos. The smallness of the neutrino Yukawa couplings is related to the presence of heavy fields whose decay provides the seed for the baryon number of our universe. The only ingredient used in this scenario above and beyond the usual leptogenesis scenario is the imposition of a $U(1)_N$ symmetry. It would be interesting to search for a fundamental origin for this symmetry. Because of the simplicity of this model, we believe that leptogenesis with Dirac neutrinos should be placed on an equal footing with the usual Majorana leptogenesis scenarios.

This model clearly displays that neutrinos need not be Majorana in order for them to play a major role in the generation of the baryon asymmetry. In this scenario, leptogenesis will not give rise to any signal in neutrinoless double beta decay experiments.

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