Warped Leptogenesis with Dirac Neutrino Masses

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Abstract

We show how leptogenesis can occur at the TeV scale with neutrinos that possess almost purely Dirac masses and negligible Majorana mass contributions as a consequence of the small wavefunction overlap in a warped fifth dimension. Lepton number violation at the Planck scale is introduced via a Majorana mass term on the Planck brane. Such a Majorana mass operator leads to the small mass splitting of otherwise degenerate Kaluza-Klein excited states on the TeV brane. This tiny mass splitting can compensate for the small Yukawa couplings to give a CP asymmetry large enough to produce the sufficient baryon asymmetry from the decay of the nearly degenerate neutrino Kaluza-Klein states. In this way the standard baryogenesis via leptogenesis scenario can naturally occur at the TeV scale without the need for a high mass scale.

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1 Introduction

The generation of a baryon asymmetry of the Universe is a fundamental question in particle physics and cosmology. The usual baryogenesis via leptogenesis scenario \[1\] occurs when a right-handed neutrino decays out of equilibrium and then sphalerons reprocess the lepton asymmetry into a baryon asymmetry above the electroweak scale. This is a theoretically appealing possibility, except that the leptogenesis scenarios involving heavy right-handed neutrinos needed for the see-saw mechanism \[2\] are difficult to experimentally verify. In addition to the difficulties in probing the nature of heavy right-handed neutrinos via energetically accessible light (left-handed) neutrinos, keeping track of the history of the universe from such a high energy scale, (e.g. GUT scale), down to the present time is not a trivial issue and it would be desirable if the observed baryon asymmetry was produced at a much lower energy scale where it could be directly accessible by terrestrial or astrophysical experiments.

We will show that the usual leptogenesis scenario can in fact be implemented at the TeV scale with predominantly Dirac neutrino masses and a small lepton number violating Majorana contribution. All small parameters naturally arise from the warped geometry. In particular the tiny Dirac Yukawa couplings result from the small wavefunction overlap of the fields in the warped fifth dimension. Interestingly, even though the neutrino is essentially Dirac in nature, there can still be sufficient lepton number violating effects to realize the leptogenesis. This is done by including a Majorana mass term on the Planck brane which represents the expected breaking of global symmetry by the higher dimensional operators induced from Planck scale physics. The right-handed neutrinos are localized near the TeV brane, so that the effects of the UV-localized Majorana mass operator on the standard model neutrino masses, are highly suppressed at the TeV brane realizing the Dirac neutrino masses. The Standard Model matter fields are localized throughout the bulk to obtain the necessary Yukawa couplings via wavefunction overlap with a Higgs field near the TeV brane \[3\].

Since the CP asymmetry due to the right-handed neutrino decay is proportional to the Yukawa couplings squared, it may seem that the resultant baryon asymmetry is too small to account for the current baryon asymmetry of the universe. However, as we shall point out, the baryon asymmetry is also inversely proportional to the small mass difference of the almost degenerate Kaluza-Klein Majorana states. Consequently, the Majorana mass term localized on the Planck brane gives a desirable tiny mass splitting between the even and odd excited Kaluza-Klein (KK) modes to compensate for the small Yukawa couplings. Thus the requisite lepton asymmetry is generated which is then reprocessed into a baryon asymmetry by electroweak sphalerons. Furthermore, in the four-dimensional (4D) interpretation of our model, the lowest lying Kaluza-Klein states are composite and are thermally produced at the TeV scale. Hence our scenario is experimentally

\[1\] If the right handed neutrinos are instead localized near the Planck brane, the Majorana mass contributions become dominant leading to the see-saw mechanism in the warped extra dimensions \[4\]. Alternatively, if one imposes lepton number conservation even on the Planck brane, then another possibility is a Dirac neutrino mass scenario which localizes the right (left) handed neutrino near the Planck (TeV) brane \[5\].
verifiable because the Kaluza-Klein states could be directly produced at the LHC.

After introducing the setup and notation of the warped extra-dimension model in § 2, we discuss the properties of the bulk neutrinos in the presence of boundary Majorana masses in § 3. In § 4 we present analytical estimates for the parameter constraints to produce the desirable baryon asymmetry of the Universe. This includes constraints arising from the electron and electron-neutrino Yukawa couplings in the standard model that are consistent with experimental observations. Finally our discussion/conclusion is in § 5.

2 Setup

Consider the fifth dimension $y$ compactified on an orbifold $S^1/Z_2$ of radius $R$, with $-\pi R \leq y \leq \pi R$, which is bounded by two three-branes at the orbifold fixed points $y = 0, \pi R$ known as the UV (or Planck) and IR (or TeV) brane respectively. The five-dimensional (5D) Einstein’s equations for this geometry lead to [5]

$$ds^2 = e^{-2k y} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2,$$

where the AdS curvature radius is $1/k$ and 4D metric is $\eta_{\mu\nu} = \text{diag}(1, -1, -1, -1)$.

The bulk action for the 5D Dirac spinor, consisting of two two-component spinors $\Psi = (\psi, \chi)^T$, has the terms [6, 7]

$$\int d^5 x \sqrt{-g} \left[ -i \bar{\psi} \sigma^\mu \partial_\mu \psi - i \bar{\chi} \sigma^\mu \partial_\mu \chi + \frac{1}{2} (\chi \overrightarrow{\partial}_5 \psi - \bar{\psi} \overrightarrow{\partial}_5 \bar{\chi}) + m_D (\chi \psi + \bar{\psi} \bar{\chi}) + \frac{1}{2} m_M (\psi \psi + \bar{\chi} \bar{\chi} + \text{h.c.}) \right],$$

where $\overrightarrow{\partial}_5 = \overrightarrow{\partial}_y - \overrightarrow{\partial}_y$. The term $\psi \partial_5 \chi$, for example, implies that $\psi \chi$ is odd under the $Z_2$ action $y \rightarrow -y$ and, for definiteness, we assign an even parity for $\psi$ and an odd parity for $\chi$ in the rest of this section. We parametrize this bulk Dirac mass, which should be odd under $Z_2$, in terms of the step function $\epsilon(y) \equiv y/|y|

$$m_D = c \, k \epsilon(y),$$

where $c$ is a dimensionless parameter. We are also interested in the right-handed neutrino Majorana mass completely localized on the Planck brane which can be parameterized as

$$m_M = d_M \delta(y),$$

where $d_M$ is a dimensionless constant. This is motivated from the fact that global symmetries are in general expected to be broken by Planck-scale physics. Hence a Majorana mass term can well be located on the Planck brane, which leads to negligible lepton-number violation on the TeV brane.
In five dimensions fermion fields can be decomposed as
\[ \psi(x, y) = \frac{e^{2ky}}{\sqrt{2\pi R}} \sum_{n=0}^{\infty} \psi^{(n)}(x) f_+^{(n)}(y), \quad \bar{\chi}(x, y) = \frac{e^{2ky}}{\sqrt{2\pi R}} \sum_{n=1}^{\infty} \bar{\chi}^{(n)}(x) f_-^{(n)}(y), \tag{5} \]
where \(+(-)\) indicates an even (odd) parity under \(Z_2\). Because the Majorana mass is confined on the Planck brane, the bulk equations of motion have the same form as those without the boundary Majorana masses and, consequently, the solutions are
\[
\text{Re} f_+^{(n)}(y) = \frac{e^{ky/2}}{N_n} \left[ J_{|c+1/2|} \left( \frac{m_n}{ke^{-ky}} \right) - J_{|c+1/2|+1} \left( \frac{m_n}{ke^{-ky}} \right) \right],
\]
\[
\text{Re} f_-^{(n)}(y) = \frac{e^{ky/2}}{N_n} \left[ J_{|c-1/2|} \left( \frac{m_n}{ke^{-ky}} \right) - J_{|c-1/2|-1} \left( \frac{m_n}{ke^{-ky}} \right) \right],
\tag{6}
\]
with \(c < -1/2 (c > -1/2)\) for the even parity field and the normalization constants \(N_n\) obtained from
\[
\frac{1}{2\pi R} \int_{-\pi R}^{\pi R} dy \, e^{ky} f_+^{(n)*}(y) f_+^{(m)}(y) = \delta_{nm}. \tag{7}
\]
The spectrum of KK masses \(m_n\) are determined by the boundary conditions (we choose the basis such that \(d_M\) is real)
\[
\text{Re} f_+^{(n)}(0) - \frac{d_M}{2} \text{Re} f_+^{(n)}(0) = 0, \\
\text{Re} f_-^{(n)}(\pi R) = 0. \tag{8}
\]
The solutions for the imaginary parts of \(f_+^{(n)}\) can be obtained by switching the sign of the Majorana mass in the boundary condition.

The analogous procedures can be applied for the KK decomposition of the scalar field \(\Phi\) as well whose action contains
\[
\int d^5x \sqrt{-g} \left( |\partial_M \Phi|^2 - m_{\phi}^2 |\Phi|^2 \right). \tag{9}
\]
For the KK decomposition
\[
\Phi(x, y) = \frac{1}{\sqrt{2\pi R}} \sum_{n=0}^{\infty} \Phi^{(n)}(x) f_n(y), \quad \frac{1}{2\pi R} \int_{-\pi R}^{\pi R} dy \, e^{-2ky} f_n(y) f_m(y) = \delta_{mn}, \tag{10}
\]
and \(m_{\phi}^2\) containing both bulk and boundary masses parameterized by
\[
m_{\phi}^2 = ak^2 + 2bk[\delta(y) - \delta(y - \pi R)], \tag{11}\]
the zero mode solution is given by

\[ f^{(0)}(y) = \frac{e^{bky}}{N_0} , \quad \frac{1}{N_0} = \sqrt{\frac{2(b - 1)\pi kR}{e^{2(b-1)\pi kR} - 1}} . \] (12)

In the above, the boundary mass of form \( b = 2 \pm \sqrt{4 + a} \) is assumed without which zero-th KK mode would vanish [7].

The bulk mass parameters \( c > 1/2 \) \((c < 1/2)\) and \( b < 1 \) \((b > 1)\) correspond to the localizations of the zero mode wave functions around Planck (TeV) brane for Fermion and scalar fields respectively.

### 3 Right-handed neutrinos in extra dimensions

We can obtain the masses of the KK states in the presence of the Majorana mass from the boundary conditions Eq. (8). However the solutions are not in a useful form for the analytical estimation of the excited KK masses, so instead we obtain the KK mass spectrum of the right-handed neutrino using the basis \( \{ \bar{f}^{(n)} \} \) which is obtained without including the Majorana mass \(^2\). In this case, the integration of \( \{ f^{(n)} \} \) over the extra dimension receives the contributions from the boundary Majorana mass terms, which gives the mixing among the KK states. Then the diagonalization of this resultant mass matrix can give the required KK state masses as performed below \(^3\). For notational clarity, we use \( \{ f^{(n)} \} \), instead of \( \{ \bar{f}^{(n)} \} \), in the following discussions to denote the basis without the Majorana mass term.

For the Majorana mass confined on the Planck brane, the KK states for a left-handed neutrino \( \nu \) and those for a right-handed neutrino \( N \) do not mix in the mass eigenstates before the electroweak symmetry is spontaneously broken by the finite Higgs VEV. The corresponding mass terms with the vanishing Higgs VEV for KK states of \( N(x, y) = (N_+(x, y), N_-(x, y))^T \) \((+(-) \) indicates even (odd) under \( Z_2 \) orbifold symmetry\) are

\[ \frac{1}{2} \begin{pmatrix} N_+^{(0)} , N_+^{(1)} , N_-^{(1)} , \ldots \end{pmatrix} \begin{pmatrix} A_{00} & A_{01} & 0 & \ldots \\ A_{01} & A_{11} & D_{N1} & \ldots \\ 0 & D_{N1} & 0 & \ldots \\ \vdots & \vdots & \vdots & \ddots \end{pmatrix} \begin{pmatrix} N_+^{(0)} \\ N_+^{(1)} \\ N_-^{(1)} \\ \vdots \end{pmatrix} , \] (13)

where we have used the basis of right-handed neutrinos such that the mass matrix elements are real. The spectrum of KK (Dirac) masses \( D_{N\bar{m}} \) for \( N \) is determined by the boundary conditions

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\(^2\)We can treat \( \bar{f}^{(n)} \) to be real because the bulk equations of motions and the boundary conditions are identical for \( f^{(n)} \) and their conjugate \( f^{(n)*} \).

\(^3\)Note that the matrix includes all the KK states up to some UV cutoff scale.
Eq. (8) with \( d_M = 0 \) and

\[
A_{mn} = \int_{-\pi R}^{\pi R} \frac{dy}{2\pi R} m_M(y) f^{(m)}_{N^+}(y) f^{(n)}_{N^+}(y) = \frac{1}{2\pi R} d_M f^{(m)}_{N^+}(0) f^{(n)}_{N^+}(0),
\]

where higher-order contributions are at least suppressed by a factor of order \( \alpha \).

Note that there is no particular reason for the phases of \( \lambda \) expected for \( d \) in the parameter range of our interest. Due to the Majorana mass term on the Planck brane, each \( \chi_1 \) nearly degenerate Majorana states.

The couplings of \( \chi_{1,2} \) to the leptons and Higgs are obtained by substituting the Kaluza-Klein decomposition ansatz, which gives the following effective 4D Lagrangian

\[
\mathcal{L}_{\text{eff}} \supset \int_{-\pi R}^{\pi R} dy \left[ \lambda_{\alpha\beta,\nu^+} L_{\alpha^+}(x, y) \tilde{H}(x, y) N_{\beta^+}(x, y) + \lambda_{\alpha\beta,\nu^-} L_{\alpha^-}(x, y) \tilde{H}(x, y) N_{\beta^-}(x, y) + \text{h.c.} \right],
\]

where \( \lambda_{\alpha\beta,\nu^\pm} \) are the 5D neutrino Yukawa couplings with mass dimension \(-1/2\), and \( \tilde{H} = i \sigma_2 H^* \).

Note that there is no particular reason for the phases of \( \lambda_{\alpha\beta,\nu^\pm} \) to be aligned for different \( Z_2 \) parity ±

\[\text{Note that, in the absence of Majorana mass term } d_M = 0, \text{ the apparent lepton number violation from the mass term } m^{(n)}_{\chi_1,\chi_2} \text{ is canceled out by that from } m^{(n)}_{\chi_2,\chi_2} \text{, so that the total lepton number is conserved as expected for } d_M = 0.\]

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or different flavor indices \((\alpha\beta)\), which will become important when we consider the CP asymmetry. We shall hereafter omit the flavor indices for brevity unless stated otherwise.

4 Leptogenesis

We shall see how much baryon asymmetry can arise from the decay of KK excited states of right-handed neutrinos to obtain the analytical estimations for the parameter constraints.

4.1 CP asymmetry

In particular, consider the lepton asymmetry arising from the decay of the first KK excited states \(\chi_{1,2}^{(1)}\). The Yukawa coupling terms relevant for \(\chi_{1,2}^{(1)}\) decays are

\[
\sum_{n=0}^{\infty} \sum_{m=0}^{\infty} \tilde{\lambda}_{\nu_\pm}^{(1,n,m)} L_\pm^{(n)} (x) \tilde{H}^{(m)} (x) N_\pm^{(1)} (x) + \sum_{n=1}^{\infty} \sum_{m=0}^{\infty} \tilde{\lambda}_{\nu_\pm}^{(1,n,m)} L_\pm^{(n)} (x) \tilde{H}^{(m)} (x) N_-^{(1)} (x) + h.c. \tag{21}
\]

\[
= \sum_n \sum_m \lambda_{\nu_\pm}^{(1,n,m)} L_\pm^{(n)} (x) \tilde{H}^{(m)} (x) \chi_1^{(1)} (x) + \lambda_{\nu_\pm}^{(1,n,m)} L_\pm^{(n)} (x) \tilde{H}^{(m)} (x) \chi_2^{(1)} (x) + \ldots, \tag{22}
\]

where \(\tilde{\lambda}_{\nu_\pm}^{(1,n,m)}\) can be obtained from Eq. (20) by the integration over the fifth dimension

\[
\tilde{\lambda}_{\nu_\pm}^{(1,n,m)} = \frac{\lambda_{\nu_\pm}^{(1,n,m)}}{(2\pi R)^{3/2}} \int_{-\pi R}^{\pi R} dy \; f_\nu^{(n)} (y) f_H^{(m)} (y) f_{N_\pm}^{(1)} (y), \tag{23}
\]

and \(\lambda_{\nu_1,2\pm}^{(1,n,m)}\) can be obtained by substituting the mass eigenstates Eq. (20) indicating the \(\tilde{\lambda}_{\nu_\pm}^{(1,n,m)}\) decays are

\[
\text{even though one may naively expect a very small}
\]

CP asymmetry due to the tiny Yukawa coupling in our Dirac mass scenario, the small mass splitting together with the small wave function overlap partially compensates the tiny Yukawa couplings.

The CP asymmetry due to the out-of-equilibrium decay of $\chi^{(1)}_2$ can give the additional contributions of the same order as that of $\chi^{(1)}_1$ for a small mass splitting $m^{(1)}_{\chi_2} - m^{(1)}_{\chi_1} \ll m^{(1)}_{\chi_1}$ as one can see by switching $\chi_1 \leftrightarrow \chi_2$ in Eq. (24).

Note also that the estimation in Eq. (24) takes into account only the self-energy contributions and there are also the vertex contributions which could give an additional CP asymmetry. We however shall not include it in the following analytical estimation for simplicity because the vertex contributions will not become much larger than self-energy contributions in our scenario, and in fact CP asymmetry of $\chi^{(1)}_1$ from the vertex corrections and that of $\chi^{(1)}_2$ cancel out as one can see by switching $\chi_1 \leftrightarrow \chi_2$ in Eq. (24).

As a concrete example to illustrate that our scenario is a workable model, let us consider the decay of electron right-handed neutrinos. In this case, potentially stringent constraints arise from

Note that $s$, $n_\gamma$ are respectively entropy and photon density and $g_s \sim O(100)$ is the number of relativistic degrees of freedom. The sphaleron effects then convert it to the baryon asymmetry via the electroweak anomaly for the observed baryon asymmetry of the universe $Y_B \sim -\frac{1}{3} Y_L \sim O(10^{-10})$.

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the electron Yukawa coupling, electron-neutrino Yukawa couplings for the zero and first KK modes and the baryon asymmetry of the Universe. These constraints can nevertheless be satisfied by adjusting the four main free parameters in our model, namely, the bulk mass parameters for the leptons and Higgs. For example, using $\pi k R \sim 34.5$ for $k e^{-\pi k R} \sim \text{TeV}$, a bulk mass parameter choice of $(c_N, c_L, c_{eR}, b_H) \sim (-0.8, 1.75, 1.2, 1.39)$ can realize $Y_B \sim \mathcal{O}(10^{-10})$ with the electron Yukawa coupling $\lambda^{(0)}_e \sim 10^{-6}$, electron-neutrino Yukawa coupling $\lambda^{(0)}_{\nu} \sim 10^{-19}$, the dominant Yukawa coupling for the first KK state $\lambda^{(1)}_{\nu} \sim 10^{-7}$ and the mass splitting of $(m_{\chi^2}^{(1)} - m_{\chi^1}^{(1)})/m_{\chi^1}^{(1)} \sim 10^{-8} d_M$ with the Planck scale Majorana mass parameter $d_M \sim 0.1$ and the 5D electron-neutrino Yukawa coupling of order $0.1$.

Note that the Majorana mass contribution to the zero mode right-handed neutrino becomes of order $A_{00} \sim 10^{-39} d_M$, and is indeed much smaller than the Dirac mass contribution. Hence our neutrino is still essentially Dirac in nature. We also point out that introducing supersymmetry does not affect our TeV-scale leptogenesis mechanism. For instance, in the parameter range where a bulk Higgs is not quite peaked on the IR brane (such as in our concrete example), supersymmetry would be required to solve the usual gauge hierarchy problem. In fact, this supersymmetric generalization would enlarge the allowable parameter space and ease constraints on model building (for example, the extra Yukawa coupling of a second Higgs doublet introduces an additional source of CP violation).

Of course to obtain a more precise quantitative estimate of the baryon asymmetry requires one to solve the Boltzmann equations which should take into account other non-trivial effects such as flavor, wash-out effects and thermal corrections \cite{14, 16}. Nevertheless our analysis shows that one can obtain an interesting viable model of the baryogenesis via leptogenesis scenario even if the Majorana mass is Dirac-like without invoking the see-saw mechanism, and where naturally small parameters are a characteristic feature of a warped extra dimension.

5 Conclusion

We have presented a leptogenesis scenario where the neutrino masses are Dirac-like and yet the production of the baryon asymmetry is possible despite tiny neutrino Yukawa couplings. A key role was played by the small mass splitting in the nearly degenerate even and odd right-handed neutrino KK states. We emphasize that both the tiny Dirac Yukawa couplings and the small mass splittings naturally arise in the warped extra dimension.

Our model also has an interesting 4D dual interpretation via the AdS/CFT correspondence \cite{3} and references therein. The right-handed neutrino KK states are composite states in the dual gauge theory, while the left-handed neutrinos are elementary. Lepton number is a global symmetry of the gauge theory, but is explicitly broken in the elementary sector. The gauge theory

\footnote{For our parameter choice, $m_{\chi^1}^{(1)}$ can decay to other first KK states such as $L_{\nu}^{(1)}$ and consequently a more precise quantitative analysis should take account of the effects of nonzero particle masses on the phase space suppression in the decay width and also on the sphaleron processes \cite{13, 14, 15}.}
couples to the elementary sector through irrelevant operators so only a small amount of lepton-number violation appears in the gauge theory. This corresponds to the small Majorana contributions to the right-handed neutrinos. The tiny Dirac couplings result from fermionic operators in the gauge theory with large anomalous dimensions. Therefore in the dual 4D theory it is the strong dynamics that is responsible for generating the small couplings and mass splittings.

While our mechanism is simple, other possibilities could also be incorporated in a warped dimension. One of the most stringent constraints came from the condition that KK states should decay before the electroweak phase transition temperature $T_c$ is reached, after which the sphaleron effects are suppressed by a factor $e^{-T_c/T}$. This constraint however can be relaxed if the parameter tuning is possible such that $m_{\chi_2}^{(1)} - m_{\chi_1}^{(1)} \sim \Gamma_{\chi_1,2}^{(1)}$, which can lead to the resonant CP asymmetry becoming as large as of order unity $\epsilon \sim \mathcal{O}(1)$ to compensate such a suppressed sphaleron rate \cite{9,10,11}.

Alternatively, if the Yukawa couplings are small enough to prevent the equilibration between the lepton asymmetry for the left and right-handed neutrino, the mechanism analogous to “neutrino-genesis” can also be another possibility to produce the baryon asymmetry in the extra dimension scenarios \cite{17}. We also point out that even though we considered the mass splitting at tree level from the Majorana mass operator confined on the Planck brane, which can be justified from the natural lepton number violation via Planck scale physics, there can also be potential effects from radiative corrections as well \cite{11,18}.

These additional effects and features deserve further study since the warped extra dimension provides an interesting alternative framework for baryon asymmetry production mechanisms at the TeV scale.

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\footnote{In the parameter range of our interest, the radiative corrections arising from Yukawa and/or graviton couplings are suppressed by the additional small Yukawa couplings and/or small neutrino Majorana masses compared with the tree level calculations, and therefore our tree level estimates of the mass splitting of first KK states would suffice for our discussions.}
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