Inflationary cosmology with five dimensional $SO(10)$

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We discuss inflationary cosmology in a five dimensional $SO(10)$ model compactified on $S^1/(Z_2\times Z_2')$, which yields $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$ below the compactification scale. The gauge symmetry $SU(5)\times U(1)_X$ is preserved on one of the fixed points, while "flipped" $SU(5)'\times U(1)_X'$ is on the other fixed point. Inflation is associated with $U(1)_X$ breaking, and is implemented through F-term scalar potentials on the two fixed points. A brane-localized Einstein-Hilbert term allows both branes to have positive tensions during inflation. The scale of $U(1)_X$ breaking is fixed from $\delta T/T$ measurements to be around 10^{16} GeV, and the scalar spectral index $n=0.98$–0.99. The inflaton field decays into right-handed neutrinos whose subsequent out of equilibrium decay yields the observed baryon asymmetry via leptogenesis.

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I. INTRODUCTION

In a recent paper [1], we showed how supersymmetric (SUSY) inflation can be realized in five dimensional (5D) models in which the fifth dimension is compactified on the orbifold $S^1/(Z_2\times Z_2')$. Orbifold symmetry breaking in higher dimensional grand unified theories (GUTs) have recently attracted a great deal of attention because of the apparent ease with which they can circumvent two particularly pressing problems encountered in four dimensional (4D) SUSY GUTs [2]: namely, the doublet-triplet (DT) splitting problem and the problem caused by dimension five nucleon decay. The apparent reluctance of the proton decay, as shown by the recent lower limits on its lifetime [3], seems to be in broad disagreement with the predicted rates from dimension five processes in minimal SUSY $SU(5)$ and $SO(10)$ models. The existence of the orbifold dimension makes it possible to implement DT splitting and simultaneously suppress (or even eliminate) dimension five proton decay.

The inflationary scenario described in Ref. [1] was inspired by the above particle physics considerations and has some novel features. The primordial density (temperature) fluctuations are proportional to $(M/M_{\rm Planck})^2$, along the lines of the 4D model proposed in Ref. [4]. Here M refers to the scale of some symmetry breaking that is associated with inflation, and $M_{\rm Planck}=1.2\times 10^{19}$ GeV denotes the Planck scale. In an $SO(10)$ model, for example, the orbifold breaking can be used to yield the subgroup $H=SU(4)_c\times SU(2)_L\times SU(2)_R$ [5], so that inflation is associated with the breaking of H to the minimal SUSY standard model (MSSM) gauge group [6,1]. The anisotropy measurements [7] can provide a determination of M independently of any particle physics considerations. M turns out to be quite close (or equal) to the SUSY GUT scale of around 10^{16} GeV. Last but not least, the scalar spectral index of density fluctuations is very close to unity ($n=0.98-0.99$) [6,1]. The gravitational wave contribution to the quadrupole anisotropy is found to be essentially negligible. In an $SO(10)$ model with the subgroup H given above, the inflaton decays into the MSSM singlet (right-handed) neutrinos, whose out of equilibrium decay leads to the observed baryon asymmetry via leptogenesis [8,9]. As we will see, this is also the case even with a different subgroup H of $SO(10)$. Baryogenesis via leptogenesis appears to be a rather generic feature of 5D $SO(10)$ based inflationary models considered here.

Although our considerations are quite general, in this paper we focus on an example based on $SO(10)$ with subgroup $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$ obtained from $S^1/(Z_2\times Z_2')$ orbifolding. The standard $SU(5)\times U(1)_X$ is preserved on one brane, while "flipped" $SU(5)'\times U(1)_X'$ is on the other brane. All massless modes from the chiral component of the 5D vector multiplet can be easily superheavy by introducing Higgs fields in the bulk. Inflation is associated with the breaking of $U(1)_X$, followed by its decay into right-handed neutrinos, which subsequently generate a primordial lepton asymmetry. The gravitino constraint on the reheat temperature [10] imposes important constraints on the masses of the right-handed neutrinos which can be folded together with the information now available from the oscillation experiments [11].

As emphasized in Ref. [1], implementation in five dimensions of the inflationary scenario considered in Ref. [4] requires some care. Note that the five dimensional setup is the appropriate one because of the proximity of the scale of inflation and the compactification scale (both are of order M_{GUT}). The inflaton potential must be localized on the orbifold fixed points (branes), since a superpotential in the bulk is not allowed. For a vanishing bulk cosmological constant, a three space inflationary solution triggered by nonzero brane tensions (or vacuum energies) exists [1]. However, 5D Einstein equations often require that the signs of the brane tensions on the two branes are opposite, which is undesirable. As shown in Ref. [1], this problem can be circumvented by introducing a brane-localized Einstein-Hilbert term in the action. The two brane tensions are both positive during inflation, and they vanish when it ends.

The plan of our paper is as follows. In Sec. II, we review cosmology in a five dimensional setting, and discuss in par-

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ticular the transition from an inflationary to the radiation dominated epoch. In Sec. III, we discuss the orbifold breaking of $SO(10)$ to $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$. Section IV summarizes the salient features of the inflationary scenario and subsequent leptogenesis. Our conclusions are summarized in Sec. V.

II. 5D COSMOLOGY

We consider 5D space-time $x^M=(x^\mu,y)$, $\mu=0,1,2,3$, compactified on an S^1/Z_2 orbifold, and the supergravity (SUGRA) action is typically given by

$$S=\int d^4x\int_{-y_c}^{y_c}dye\left[\frac{M_5^3}{2}R_5+\mathcal{L}_B+\sum_{i=I,II}\frac{\delta(y-y_i)}{e_5^5}\times\left(\frac{M_i^2}{2}\overline{R}_4+\mathcal{L}_i\right)\right], \tag{1}$$

where R_5 ($\overline{R}_4$) stands for the 5 dimensional (4 dimensional) Einstein-Hilbert term, $\mathcal{L}_B$ ($\mathcal{L}_I$, $\mathcal{L}_{II}$) denotes some unspecified bulk (brane) contributions to the full Lagrangian, and $y_I=0$, $y_{II}=y_c$ indicate the brane positions. The brane scalar curvature term $\overline{R}_4(\overline{g}_{\mu\nu})$ is defined through the induced metric, $\overline{g}_{\mu\nu}(x)\equiv g_{\mu\nu}(x,y=0)$ $(\mu,\nu=0,1,2,3)$. The brane-localized Einstein-Hilbert terms[1] in Eq. (1) are allowed also in SUGRA, but should, of course, be accompanied by brane gravitino kinetic terms as well as other terms, as spelled out in the off-shell SUGRA formalism [13]. Here we assume that the bulk cosmological constant is zero.

For the cosmological solution let us take the following metric ansatz,

$$ds^2=\beta^2(t,y)[-dt^2+a^2(t)d\vec{x}^2]+dy^2, \tag{2}$$

which shows that the three dimensional space is homogeneous and isotropic. The nonvanishing components of the 5D Einstein tensor derived from Eq. (1) are

$$G_0^0=3\left[\left(\frac{\beta''}{\beta}\right)+\left(\frac{\beta'}{\beta}\right)^2\right]-\frac{3}{\beta^2}\left[\left(\frac{\dot\beta}{\beta}+\frac{\dot a}{a}\right)^2\right]-\sum_{i=I,II}\delta(y-y_i)\frac{M_i^2}{M_5^3}\frac{3}{\beta^2}\left[\left(\frac{\dot\beta}{\beta}+\frac{\dot a}{a}\right)^2\right], \tag{3}$$

[1]The importance of the brane-localized 4D Einstein-Hilbert term, especially for generating 4D gravity in a higher dimensional noncompact flat space was first noted in Ref. [12].

$$G_i^i=3\left[\left(\frac{\beta''}{\beta}\right)+\left(\frac{\beta'}{\beta}\right)^2\right]-\frac{1}{\beta^2}\left[2\frac{\ddot\beta}{\beta}+2\frac{\ddot a}{a}+4\frac{\dot\beta}{\beta}\frac{\dot a}{a}-\left(\frac{\dot\beta}{\beta}\right)^2+\left(\frac{\dot a}{a}\right)^2\right]-\sum_{i=I,II}\delta(y-y_i)\frac{M_i^2}{M_5^3}\frac{1}{\beta^2}\times\left[2\frac{\ddot\beta}{\beta}+2\frac{\ddot a}{a}+4\frac{\dot\beta}{\beta}\frac{\dot a}{a}-\left(\frac{\dot\beta}{\beta}\right)^2+\left(\frac{\dot a}{a}\right)^2\right], \tag{4}$$

$$G_5^5=6\left[\left(\frac{\beta'}{\beta}\right)^2\right]-\frac{3}{\beta^2}\left[\frac{\ddot\beta}{\beta}+\frac{\ddot a}{a}+3\frac{\dot\beta}{\beta}\frac{\dot a}{a}+\left(\frac{\dot a}{a}\right)^2\right], \tag{5}$$

$$G_{05}=-3\left[\left(\frac{\beta'}{\beta}\right)^{\cdot}\right], \tag{6}$$

where primes and dots respectively denote derivatives with respect to y and t, and the terms accompanied by delta functions arise from the brane-localized Einstein-Hilbert terms.

Let us first discuss inflation under this setup. Since 5D $N=1$ SUSY does not allow a superpotential (and the corresponding F-term scalar potential) in the bulk, we introduce the inflaton scalar potentials $V_{I,II}(\phi)$ $(\geqslant 0)$ on the two branes where only 4D $N=1$ SUSY is preserved [1,14]. The energy-momentum tensor during inflation is given by

$$T_0^0=T_i^i=-\delta(y)\frac{V_I}{M_5^3}-\delta(y-y_c)\frac{V_{II}}{M_5^3}, \tag{7}$$

$$T_5^5=0, \tag{8}$$

where $i=1,2,3$, and V_I, V_{II} are the scalar potentials on the branes at $y=0$ (B1) and $y=y_c$ (B2), respectively. They are suitably chosen to provide a large enough number of e-foldings to resolve the horizon and flatness problems. The end of inflation is marked by the breaking of the "slow roll" conditions, and the inflaton rolls quickly to the true suepersymmetric vacuum with flat 4D space-time. Thus, for the inflationary epoch it is sufficient to consider only scalar potentials in the energy-momentum tensor. We will discuss more general cases later.

The exact inflationary solution is [1]

$$\beta(y)=H_0|y|+c, \tag{9}$$

$$a(t)=e^{H_0t}, \tag{10}$$

where H_0 (>0) is the Hubble constant during inflation. The integration constant c in Eq. (9) can be normalized to unity without loss of generality. The introduction of the brane-localized Einstein-Hilbert terms do not affect the bulk solutions, Eqs. (9) and (10), but they modify the boundary conditions. The solution $\beta(y)$ should satisfy the following boundary conditions at $y=0$ and $y=y_c$,

$$\frac{V_I}{6M_5^3}=-H_0+\frac{1}{2}\frac{M_I^2}{M_5^3}H_0^2, \tag{11}$$

$$\frac{V_{II}}{6M_5^3}=\frac{H_0}{1+H_0y_c}+\frac{1}{2}\frac{M_{II}^2}{M_5^3}\frac{H_0^2}{(1+H_0y_c)^2}. \tag{12}$$

Thus, H_0 and y_c are determined by V_I and V_{II}. Note that the brane cosmological constants (scalar potentials) V_I and V_{II} are related to the Hubble constant H_0. While the nonzero brane cosmological constants are responsible for inflating the 3-space, their subsequent vanishing restores SUSY and guarantees the flat 4D space-time. Since V_{II} must be zero when $V_I=0$, it is natural that the scalar field controlling the end of inflation is introduced in the bulk. With SUSY broken at low energies, the minima of the inflaton potentials on both branes should be fine-tuned to zero [15,16].

From Eqs. (11) and (12) we note that in the absence of the brane-localized Einstein-Hilbert term at $y=0$, the inflaton potentials (brane cosmological constants) V_I and V_{II} should have opposite signs. However, a suitably large value of M_I/M_5 [17] can flip the sign of V_I [1], so that both V_I and V_{II} are positive. Thus, a brane-localized Einstein-Hilbert term at $y=0$ seems essential for successful F-term inflation in the 5D SUSY framework. Its introduction does not conflict with any symmetry, and in Ref. [1] a simple model for realizing a large ratio M_I/M_5 was proposed.

The 4D reduced Planck mass $[\equiv(M_{\rm Planck}/8\pi)^{1/2}]$ is given by

$$\begin{aligned}M_P^2&=M_5^3\int_{-y_c}^{y_c}dy\,\beta^2+M_I^2+M_{II}^2\\&=M_5^3y_c\left(\frac{2}{3}H_0^2y_c^2+2H_0y_c+2\right)+M_I^2+M_{II}^2,\end{aligned} \tag{13}$$

while the 4D effective cosmological constant is calculated to be

$$\begin{aligned}\Lambda_{\rm eff}&=\int_{-y_c}^{y_c}dy\,\beta^4\left\{M_5^3\left[4\left(\frac{\beta''}{\beta}\right)+6\left(\frac{\beta'}{\beta}\right)^2\right]\right.\\&\qquad\left.+\delta(y)V_I+\delta(y-y_c)V_{II}\right\}\\&=3H_0^2\left[M_5^3y_c\left(\frac{2}{3}H_0^2y_c^2+2H_0y_c+2\right)+M_I^2+M_{II}^2\right]\\&=3H_0^2M_P^2,\end{aligned} \tag{14}$$

which vanishes when $V_I=V_{II}=0$.

After inflation, the inflaton decays into brane and (subsequently) bulk fields, which reheat the whole 5 dimensional universe. To quantify the inflaton and radiation (or matter) dominated epochs, we use the fluid approximation,

$$T^M{}_N=\frac{1}{M_5^3}\begin{pmatrix}-\rho&0&0&0&T^0{}_5\\0&p&0&0&0\\0&0&p&0&0\\0&0&0&p&0\\T^5{}_0&0&0&0&P_5\end{pmatrix}, \tag{15}$$

where ρ and p are contributed by bulk and brane matter,

$$\rho\equiv\frac{1}{2y_c}\rho_B+\sum_{i=I,II}\delta(y-y_i)\rho_i, \tag{16}$$

$$p\equiv\frac{1}{2y_c}p_B+2y_c\sum_{i=I,II}\delta(y-y_i)p_i. \tag{17}$$

Note that in Eq. (15) the nonzero off-diagonal components, $T^0{}_5\,[=(-1/\beta^2)T_{05}]$ and $T^5{}_0\ (=T_{05})$ are considered. In Eqs. (16) and (17), we normalize ρ_B and p_B with the circumference of the extra dimension, so their components have the same mass dimension as their brane counterparts. With Eqs. (3)–(6) and (15), the 5D "Friedmann-like" equations are readily written,

$$\frac{1}{2y_cM_5^3}\rho_B=\frac{3}{\beta^2}\left[\left(\frac{\dot\beta}{\beta}\right)^2+2\frac{\dot\beta}{\beta}H+(H^2-h^2)\right], \tag{18}$$

$$\frac{1}{M_5^3}\rho_I=\left[3\frac{M_I^2}{M_5^3}H^2-6h\right], \tag{19}$$

$$\frac{1}{M_5^3}\rho_{II}=\left\{\frac{M_{II}^2}{M_5^3}\frac{3}{\beta^2}\left[\left(\frac{\dot\beta}{\beta}\right)^2+2\frac{\dot\beta}{\beta}H+H^2\right]+6\frac{h}{\beta}\right\}_{y=y_c}, \tag{20}$$

$$\frac{1}{2y_cM_5^3}p_B=-\frac{1}{\beta^2}\left[2\frac{\ddot\beta}{\beta}-\left(\frac{\dot\beta}{\beta}\right)^2+4\frac{\dot\beta}{\beta}H+2\dot H+3(H^2-h^2)\right], \tag{21}$$

$$\frac{1}{M_5^3}p_I=-\left[\frac{M_I^2}{M_5^3}(2\dot H+3H^2)-6h\right], \tag{22}$$

$$\begin{aligned}\frac{1}{M_5^3}p_{II}=-\left\{\frac{M_{II}^2}{M_5^3}\frac{1}{\beta^2}\left[2\frac{\ddot\beta}{\beta}-\left(\frac{\dot\beta}{\beta}\right)^2+4\frac{\dot\beta}{\beta}H+2\dot H+3H^2\right]\right.\\\left.+6\frac{h}{\beta}\right\}_{y=y_c},\end{aligned} \tag{23}$$

$$\frac{1}{2y_cM_5^3}P_5=-\frac{3}{\beta^2}\left[\frac{\ddot\beta}{\beta}+3\frac{\dot\beta}{\beta}H+\dot H+2(H^2-h^2)\right], \tag{24}$$

$$\frac{1}{2y_c M_5^3}T_{05} = -3\,\mathrm{sgn}(y)\left[\frac{\dot{h}}{\beta} - \frac{h}{\beta}\frac{\dot{\beta}}{\beta}\right]. \tag{25}$$

Here $\mathrm{sgn}(y)\equiv 1\ (-1)$ for $y>0\ (<0)$, and

$$H(t)\equiv\frac{\dot{a}}{a}, \tag{26}$$

$$\beta(t,y) = h(t)|y| + 1. \tag{27}$$

For $M_I > M_5, M_{II}$, and $H \gg h$, the brane matter contribution from B1 is dominant, and Eqs. (18)–(25) reduce to the approximate 4D Friedmann equations,

$$\left(\frac{\dot{a}}{a}\right)^2 \approx \frac{1}{3M_4^2}\rho_I, \tag{28}$$

$$\frac{\ddot{a}}{a} \approx \frac{-1}{6M_4^2}(\rho_I + 3p_I). \tag{29}$$

Equations (18)–(25) satisfy the energy-momentum conservation law, $\nabla_M T^M_N = 0$ whose $N=0$ and $N=5$ components are [18]

$$\dot{\rho} + 3\left(\frac{\dot{\beta}}{\beta} + H\right)(\rho + p) = T^{5\prime}{}_0 + 4\frac{\beta'}{\beta}T^5{}_0 = 2y_c M_5^3\left[G'_{05} + 4\frac{\beta'}{\beta}G_{05}\right], \tag{30}$$

$$P'_5 + \frac{\beta'}{\beta}(4P_5 - 3p + \rho) = -\dot{T}^0{}_5 - \left(4\frac{\dot{\beta}}{\beta} + 3H\right)T^0{}_5 = 2y_c M_5^3\frac{1}{\beta^2}\left[\dot{G}_{05} + \left(2\frac{\dot{\beta}}{\beta} + 3H\right)G_{05}\right]. \tag{31}$$

The inflaton contributes to the energy momentum tensor, Eq. (15),

$$T_{MN} \equiv T^{\mathrm{inf}}_{MN} + T^{\mathrm{m}}_{MN}, \tag{32}$$

where T^{inf}_{MN} denotes the contributions to the energy momentum tensor from the inflaton $\phi(t,y)$,

$$T^{\mathrm{inf}}_{MN} \equiv \frac{1}{2y_c}\partial_M\phi\partial_N\phi - \frac{1}{4y_c}g_{MN}\partial^P\phi\partial_P\phi + \sum_{i=I,II}\delta(y-y_i) \times \delta^\mu_M\delta^\nu_N\left[\partial_\mu\phi\partial_\nu\phi - g_{\mu\nu}\left(\frac{1}{2}\partial^\lambda\phi\partial_\lambda\phi + V_i(\phi)\right)\right], \tag{33}$$

and T^{m}_{MN} is assumed to have the same form as Eq. (15). The conservation law $\nabla^M T^{\mathrm{inf}}_{MN} = 0$ gives rise to the scalar field equation in the presence of both the brane and bulk kinetic terms. If only the inflaton potentials on the branes, $V_I(\phi)$ and $V_{II}(\phi)$ are dominant in Eq. (32), one can check that the solutions reduce to Eqs. (9) and (10), namely, $H = h = \mathrm{constant}\ (=H_0)$. The inflaton decay produces T^{m}_{MN}.

We have tacitly assumed that the interval separating the two branes (orbifold fixed points) remains fixed during inflation. The dynamics of the orbifold fixed points, unlike the D-brane case [19], is governed only by the $g_{55}(x,y)$ component of the metric tensor. The real fields e^5_5, B_5, and the chiral fermion ψ^2_{5R} in 5D gravity multiplet are assigned even parity under Z_2 [13], and they compose an $N=1$ chiral multiplet on the branes. The associated superfield can acquire a superheavy mass and its scalar component can develop a vacuum expectation value (VEV) on the brane. With superheavy brane-localized mass terms, the low-lying Kaluza-Klein (KK) mass spectrum is shifted so that even the lightest mode obtains a compactification scale mass [20]. Since this mass is much greater than H_0 the interval distance is stable even during inflation. This stabilization of the interval distance in turn leads to the stabilization of the warp factor $\beta(y)$. This is because the fluctuation $\delta\beta(y)$ of the warp factor near the solution in Eq. (9) turns out to be proportional to the interval length variation δg_{55} from the linearized 5D Einstein equation [21].

So far we have discussed only S^1/Z_2 orbifold compactification. The results can be directly extended to $S^1/(Z_2 \times Z'_2)$. Within the framework discussed in this section, we can accommodate any promising 4D SUSY inflationary model. We consider one particular model below which comes from compactifying $SO(10)$ on an $S^1/(Z_2\times Z'_2)$.

III. $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$ MODEL

We consider the $N=2$ (in the 4D sense) SUSY $SO(10)$ model in 5D space-time, where the 5th dimension is compactified on an $S^1/(Z_2\times Z'_2)$. Z_2 reflects $y\to -y$, and Z'_2 reflects $y'\to -y'$, with $y' = y + y_c/2$. There are two independent orbifold fixed points (branes) at $y=0$ and $y=y_c/2$. The $S^1/(Z_2\times Z'_2)$ orbifold compactification is exploited to yield $N=1$ SUSY as well as break $SO(10)$ to some suitable subgroup.

Under $SU(5)\times U(1)_X$, the $SO(10)$ generators are split into[2]

$$T_{SO(10)} = \begin{bmatrix} \mathbf{24}_0 + \mathbf{1}_0 & \mathbf{10}_{-4} \\ \overline{\mathbf{10}}_4 & \overline{\mathbf{24}}_0 - \mathbf{1}_0 \end{bmatrix}_{10\times 10}, \tag{34}$$

where the subscripts labeling the $SU(5)$ representations indicate $U(1)_X$ charges, and the subscript "10×10" denotes the matrix dimension. Also, $\mathbf{24}(=\overline{\mathbf{24}})$ corresponds to $SU(5)$

[2]The $SO(2n)$ generators are represented as $\begin{pmatrix} A+C & B+S \\ B-S & A-C \end{pmatrix}$, where A, B, C are $n\times n$ antisymmetric matrices and S is an $n\times n$ symmetric matrix [22]. By a unitary transformation, the generators are given by $\begin{pmatrix} A-iS & C+iB \\ C-iB & A+iS \end{pmatrix}$, where A and S denote $U(n)$ generators, and $C\pm iB$ transform as $n(n-1)/2$ and $\overline{n(n-1)/2}$ under $SU(n)$.

generators while $\mathrm{diag}(\mathbf{1}_{5\times 5},-\mathbf{1}_{5\times 5})$ is the $U(1)_X$ generator. The 5×5 matrices $\mathbf{24}_0$ and $\mathbf{10}_{-4}$ are further decomposed under $SU(3)_c\times SU(2)_L\times U(1)_Y$ as

$$\mathbf{24}_0=\begin{pmatrix} (\mathbf{8},\mathbf{1})_0+(\mathbf{1},\mathbf{1})_0 & (\mathbf{3},\overline{\mathbf{2}})_{-5/6} \\ (\overline{\mathbf{3}},\mathbf{2})_{5/6} & (\mathbf{1},\mathbf{3})_0-(\mathbf{1},\mathbf{1})_0 \end{pmatrix}_0, \tag{35}$$

$$\mathbf{10}_{-4}=\begin{pmatrix} (\overline{\mathbf{3}},\mathbf{1})_{-2/3} & (\mathbf{3},\mathbf{2})_{1/6} \\ (\mathbf{3},\mathbf{2})_{1/6} & (\mathbf{1},\mathbf{1})_1 \end{pmatrix}_{-4}. \tag{36}$$

Thus, each representation carries two independent $U(1)$ charges. Note that the two $(\mathbf{3},\mathbf{2})_{1/6}$s in $\mathbf{10}_{-4}$ are identified.

Consider the two independent Z_2 and Z_2' group elements,

$$P\equiv \mathrm{diag}(I_{3\times 3},\ I_{2\times 2},-I_{3\times 3},-I_{2\times 2}), \tag{37}$$

$$P'\equiv \mathrm{diag}(-I_{3\times 3},\ I_{2\times 2},\ I_{3\times 3},-I_{2\times 2}), \tag{38}$$

which satisfy $P^2=P'^2=\mathbf{1}_{5\times 5}$. Under the operations $PT_{SO(10)}P^{-1}$ and $P'T_{SO(10)}P'^{-1}$, the matrix elements of $T_{SO(10)}$ transform as

$$\begin{bmatrix} (\mathbf{8},\mathbf{1})_0^{++} & (\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{+-} & (\overline{\mathbf{3}},\mathbf{1})_{-2/3}^{--} & (\mathbf{3},\mathbf{2})_{1/6}^{-+} \\ (\overline{\mathbf{3}},\mathbf{2})_{5/6}^{+-} & (\mathbf{1},\mathbf{3})_0^{++} & (\mathbf{3},\mathbf{2})_{1/6}^{-+} & (\mathbf{1},\mathbf{1})_1^{--} \\ (\mathbf{3},\mathbf{1})_{2/3}^{--} & (\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{-+} & (\mathbf{8},\mathbf{1})_0^{++} & (\overline{\mathbf{3}},\mathbf{2})_{5/6}^{+-} \\ (\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{-+} & (\mathbf{1},\mathbf{1})_{-1}^{--} & (\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{+-} & (\mathbf{1},\mathbf{3})_0^{++} \end{bmatrix}_{10\times 10}, \tag{39}$$

where the superscripts of the matrix elements indicate the eigenvalues of the P and P' operations. Here we omitted the two $U(1)$ generators [$(\mathbf{1},\mathbf{1})_\mathbf{0}^{\mathbf{++}}$s] to avoid too much clutter. As shown in Eqs. (34) and (35), they should appear in the diagonal part of the matrix (39). The eigenvalues of P and P' are the imposed parities (or boundary conditions) of fields in the adjoint representations. The wave function of a field with parity $(+-)$, for instance, must vanish on the brane at $y=y_c/2$, and only fields assigned $(++)$ parities contain massless modes in their KK spectrum.

An $N=2$ gauge multiplet consists of an $N=1$ vector multiplet $[V^a=(A_\mu^a,\lambda^{1a})]$ and an $N=1$ chiral multiplet $[\Sigma^a=((\Phi^a+iA_5^a)/\sqrt{2},\lambda^{2a})]$. In order to break $N=2$ SUSY to $N=1$, opposite parities must be assigned to the vector and the chiral multiplets in the same representation. The parities of $N=1$ vector multiplets coincide with the parities of the corresponding $SO(10)$ generators in Eq. (39). From the assigned eigenvalues in Eq. (39), only the gauge multiplets associated with $(\mathbf{8},\mathbf{1})_0^{++}$, $(\mathbf{1},\mathbf{3})_0^{++}$, and two $(\mathbf{1},\mathbf{1})_0^{++}$, which correspond to the $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$ generators, contain massless modes. Therefore, at low energy the theory is effectively described by a $4D$ $N=1$ supersymmetric theory with $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$.

As seen in Eq. (39), all of the gauge multiplets associated with the diagonal components $\mathbf{24}_0$, $\mathbf{1}_0$ in Eq. (34) survive at B1. Thus, on B1 $SU(5)\times U(1)_X$ is preserved [23]. On the other hand, in Eq. (39), the elements with $(++)$ and $(-+)$ parities can compose a second (distinct) set of $SU(5)'\times U(1)_X'$ gauge multiplets. In the $SU(5)$ generator at B1, $\mathbf{24}_0$ $[=(\mathbf{8},\mathbf{1})_0^{++}+(\mathbf{1},\mathbf{3})_0^{++}+(\mathbf{1},\mathbf{1})_0^{++}+(\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{+-}+(\overline{\mathbf{3}},\mathbf{2})_{5/6}^{+-}]$, the $(\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{+-}$ and $(\overline{\mathbf{3}},\mathbf{2})_{5/6}^{+-}$ are replaced by $(\mathbf{3},\mathbf{2})_{1/6}^{-+}$ and $(\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{-+}$ at B2 which belong in $\mathbf{10}_{-4}$ and $\overline{\mathbf{10}}_4$ respectively of $SU(5)\times U(1)_X$,

$$\mathbf{24}_0'=(\mathbf{8},\mathbf{1})_0^{++}+(\mathbf{1},\mathbf{3})_0^{++}+(\mathbf{1},\mathbf{1})_0^{++}+(\mathbf{3},\mathbf{2})_{1/6}^{-+} +(\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{-+} \quad \text{at B2}, \tag{40}$$

where the assigned hypercharges coincide with those given in "flipped" $SU(5)'\times U(1)_X'$ [24]. The $U(1)_X'$ generator at B2 is defined as

$$\mathrm{diag}(\mathbf{1}_{3\times 3},-\mathbf{1}_{2\times 2},-\mathbf{1}_{3\times 3},\mathbf{1}_{2\times 2}). \tag{41}$$

Thus, the $U(1)_X'$ charges of the surviving elements at B2 turn out to be zero, while the other components are assigned -4 or 4. The $U(1)_X'$ generator and the matrix elements with $(++)$, $(-+)$ parities in Eq. (39) can be block-diagonalized to the form in Eq. (34) through the unitary transformation

$$U=\begin{pmatrix} I_{3\times 3} & 0 & 0 & 0 \\ 0 & 0 & 0 & I_{2\times 2} \\ 0 & 0 & I_{3\times 3} & 0 \\ 0 & I_{2\times 2} & 0 & 0 \end{pmatrix}_{10\times 10}. \tag{42}$$

We conclude that the gauge multiplets surviving at B2 are associated with a second ("flipped") $SU(5)'\times U(1)_X'$ embedded in $SO(10)$ [24].

With opposite parities assigned to the chiral multiplets, two vectorlike pairs $\Sigma_{(\overline{\mathbf{3}},\mathbf{1})_{-2/3}^{++}}$, $\Sigma_{(\mathbf{3},\mathbf{1})_{2/3}^{++}}$ and $\Sigma_{(\mathbf{1},\mathbf{1})_{1}^{++}}$, $\Sigma_{(\mathbf{1},\mathbf{1})_{-1}^{++}}$ contain massless modes. Although the nonvanishing chiral multiplets at B1 are $\Sigma_{\mathbf{10}_{-4}}$ $(=\Sigma_{(\overline{\mathbf{3}},\mathbf{1})_{-2/3}^{++}}+\Sigma_{(\mathbf{3},\mathbf{2})_{1/6}^{+-}}+\Sigma_{(\mathbf{1},\mathbf{1})_{1}^{++}})$ and $\Sigma_{\overline{\mathbf{10}}_4}$ $(=\Sigma_{(\mathbf{3},\mathbf{1})_{2/3}^{++}}+\Sigma_{(\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{+-}}+\Sigma_{(\mathbf{1},\mathbf{1})_{-1}^{++}})$, $\Sigma_{(\mathbf{3},\mathbf{2})_{1/6}^{+-}}$ and $\Sigma_{(\overline{\mathbf{3}},\overline{\mathbf{2}})_{-1/6}^{+-}}$ are replaced by $\Sigma_{(\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{-+}}$ and $\Sigma_{(\overline{\mathbf{3}},\mathbf{2})_{5/6}^{-+}}$ at B2 that are contained in $\Sigma_{\mathbf{24}_0}$ and $\Sigma_{\overline{\mathbf{24}}_0}$ at B1. Together with the vectorlike pairs containing massless modes, they compose $\mathbf{10}'_{-4}$- and $\overline{\mathbf{10}}_\mathbf{4}'$-plets of $SU(5)'\times U(1)_X'$,

$$\Sigma_{\mathbf{10}'_{-4}}=\Sigma_{(\overline{\mathbf{3}},\mathbf{1})_{-2/3}^{++}}+\Sigma_{(\mathbf{3},\overline{\mathbf{2}})_{-5/6}^{-+}}+\Sigma_{(\mathbf{1},\mathbf{1})_{-1}^{++}}, \tag{43}$$

$$\Sigma_{\overline{\mathbf{10}}_4'}=\Sigma_{(\mathbf{3},\mathbf{1})_{2/3}^{++}}+\Sigma_{(\overline{\mathbf{3}},\mathbf{2})_{5/6}^{-+}}+\Sigma_{(\mathbf{1},\mathbf{1})_{1}^{++}} \quad \text{at B2}. \tag{44}$$

TABLE I. $U(1)_R$ charges of the vector and hypermultiplets.

$U(1)_R$	V, Σ			H, H^c
1				ϕ^c
1/2		λ^1		ψ^c
0	A_μ		Φ, A_5	ϕ
$-1/2$		λ^2		ψ

TABLE II. $U(1)_X^{(\prime)}$, $U(1)_R$, $U(1)_{PQ}$ charges and matter parities of the superfields.

Fields	S	N_H	$\overline{N}_H$	$\mathbf{10}_B^{(\prime)}$	$\overline{\mathbf{10}}_B^{(\prime)}$	$\mathbf{1}_i$	$\overline{\mathbf{5}}_i$	$\mathbf{10}_i$
$X^{(\prime)}$	0	5	-5	-4	4	5	-3	1
R	1	0	0	0	0	1/2	1/2	1/2
PQ	0	0	0	0	0	0	-1	$-1/2$
Z_2^m	$+$	$+$	$+$	$-$	$-$	$-$	$-$	$-$
Fields	$h_1^{(\prime)}$	$h_2^{(\prime)}$	$\overline{h}_1^{(\prime)}$	$\overline{h}_2^{(\prime)}$	Σ_1	Σ_2	$\overline{\Sigma}_1$	$\overline{\Sigma}_2$
$X^{(\prime)}$	-2	-2	2	2	0	0	0	0
R	0	1	0	1	1/2	1/2	0	0
PQ	1	1	3/2	3/2	-1	$-3/2$	1	3/2
Z_2^m	$+$	$-$	$+$	$-$	$+$	$+$	$+$	$+$

Now let us discuss the $N=2$ (bulk) hypermultiplet H $[=(\phi,\psi)]$, $H^c[=(\phi^c,\psi^c)]$ in the vector representations $\mathbf{10}$, $\mathbf{10^c}$ $(=\mathbf{10})$ of $SO(10)$, where H and H^c are $N=1$ chiral multiplets. Under $SU(5)\times U(1)_X$ and $SU(3)_c\times SU(2)_L\times U(1)_Y$, $\mathbf{10}$ and $\mathbf{10^c}$ are

$$\mathbf{10}=\begin{pmatrix}\mathbf{5}_{-2}\\ \overline{\mathbf{5}}_2\end{pmatrix}=\begin{pmatrix}(\mathbf{3},\mathbf{1})_{-1/3}^{+-}\\ (\mathbf{1},\mathbf{2})_{1/2}^{++}\\ (\overline{\mathbf{3}},\mathbf{1})_{1/3}^{-+}\\ (\mathbf{1},\overline{\mathbf{2}})_{-1/2}^{--}\end{pmatrix},$$

$$\mathbf{10^c}=\begin{pmatrix}\mathbf{5^c}_2\\ \overline{\mathbf{5^c}}_{-2}\end{pmatrix}=\begin{pmatrix}(\mathbf{3^c},\mathbf{1})_{1/3}^{-+}\\ (\mathbf{1},\mathbf{2^c})_{-1/2}^{--}\\ (\overline{\mathbf{3^c}},\mathbf{1})_{-1/3}^{+-}\\ (\mathbf{1},\overline{\mathbf{2^c}})_{1/2}^{++}\end{pmatrix}, \tag{45}$$

where the subscripts ± 2 are $U(1)_X$ charges and the remaining subscripts indicate the hypercharges Y. The superscripts on the matrix elements denote the eigenvalues of the P and P' operations. As in the $N=2$ vector multiplet, opposite parities must be assigned for H and H^c to break $N=2$ SUSY to $N=1$. The massless modes are contained in the two doublets $(\mathbf{1},\mathbf{2})_{1/2}^{++}$ and $(\mathbf{1},\overline{\mathbf{2^c}})_{1/2}^{++}$ $[=(\mathbf{1},\mathbf{2})_{1/2}^{++}]$. While the surviving representations at B1, $(\mathbf{3},\mathbf{1})_{-1/3}^{+-}$ and $(\mathbf{1},\mathbf{2})_{1/2}^{++}$ [also $(\overline{\mathbf{3^c}},\mathbf{1})_{-1/3}^{+-}$ and $(\mathbf{1},\overline{\mathbf{2^c}})_{1/2}^{++}$] compose two $\mathbf{5}_{-2}$ (or $\overline{\mathbf{5^c}}_{-2}$) of $SU(5)\times U(1)_X$, at B2 the nonvanishing representations are two $\overline{\mathbf{5}}_2'$ of $SU(5)'\times U(1)_X'$,

$$\overline{\mathbf{5}}_2'=(\overline{\mathbf{3}},\mathbf{1})_{1/3}^{-+}+(\mathbf{1},\mathbf{2})_{1/2}^{++}\quad [\text{or}\quad (\mathbf{3^c},\mathbf{1})_{1/3}^{-+}+(\mathbf{1},\overline{\mathbf{2^c}})_{1/2}^{++}]\quad \text{at B2}. \tag{46}$$

In this model, the $SU(2)$ R symmetry which generally exists in $N=2$ supersymmetric theories is explicitly broken to $U(1)_R$. Since $N=1$ SUSY is present on both branes, $U(1)_R$ symmetry should be respected. We note that different $U(1)_R$ charges can be assigned to $H_{\mathbf{10}_{-4}}$ and $H^c_{\mathbf{10^c}_4}$ as shown in Table I [25].

To construct a realistic model, which includes inflation, based on $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$, we introduce a $U(1)_{PQ}$ axion symmetry and "matter" parity Z_2^m [6]. For simplicity, let us assume that the MSSM matter superfields as well as the right-handed neutrinos are brane fields residing at B1.[3] They belong in $\mathbf{10}_i$, $\overline{\mathbf{5}}_i$, and $\mathbf{1}_i$ of $SU(5)$, where i is the family index. Their assigned $U(1)_X$, $U(1)_R$ and $U(1)_{PQ}$ charges and matter parities appear in Table II.

We introduce two pairs of hypermultiplets $(H_{\mathbf{10}},H^c_{\mathbf{10^c}})$ and $(H_{\overline{10}},H^c_{\overline{\mathbf{10^c}}})$ $[=(H_{\mathbf{10}},H^c_{\mathbf{10^c}})]$ in the bulk. The two $SU(5)$ Higgs multiplets h_1 and $\overline{h}_1$ ($\mathbf{5}$ and $\overline{\mathbf{5}}$) arise from $H_{\mathbf{10}}$ and $H_{\overline{10}}$, and their $U(1)_R$ charges are chosen to be zero. As discussed above, the $N=2$ superpartners $H^c_{\mathbf{10}}$ and $H^c_{\overline{10}}$ also provide superfields h_2 and $\overline{h}_2$ with $\mathbf{5}^{(\prime)}$ and $\overline{\mathbf{5}}^{(\prime)}$ representations at B1 (B2). However, their $U(1)_R$ charges are unity unlike h_1 and $\overline{h}_1$. To make them superheavy we can introduce another pair of $\mathbf{5}$ and $\overline{\mathbf{5}}$ with zero $U(1)_R$ charges and $-$ matter parities on the brane.

The superpotential at B1, neglecting the superheavy particles' contributions except for the inflatons, is given by

$$\begin{aligned}W=&\kappa S(N_H\overline{N}_H-M^2)+\frac{\sigma_1}{M_P}\Sigma_1\Sigma_2h_1\overline{h}_1+\frac{\sigma_2}{M_P}\Sigma_1\Sigma_2\overline{\Sigma}_1\overline{\Sigma}_2\\&+y_{ij}^{(d)}\mathbf{10}_i\mathbf{10}_jh_1+y_{ij}^{(ul)}\mathbf{10}_i\overline{\mathbf{5}}_j\overline{h}_1+y_{ij}^{(n)}\mathbf{1}_i\overline{\mathbf{5}}_jh_1\\&+\frac{y_{ij}^{(m)}}{M_P}\mathbf{1}_i\mathbf{1}_j\overline{N}_H\overline{N}_H,\end{aligned} \tag{47}$$

where S, N_H, $\overline{N}_H$, $\Sigma_{1,2}$, and $\overline{\Sigma}_{1,2}$ are singlet fields. Their assigned quantum numbers appear in Table II. While S, N_H, $\overline{N}_H$, and h_1, $\overline{h}_1$ are bulk fields, the rest are brane fields

[3]If the first two quark and lepton families reside on B2 where $SU(5)'\times U(1)_X'$ is preserved, undesirable mass relations between the down-type quarks and the charged leptons do not arise. Mixings between the first two and the third families can be generated by introducing bulk superheavy hypermultiplets in the spinor representations of $SO(10)$ [26].

residing on B1. N_H, $\overline{N}_H$ should be embedded in $\mathbf{16}_H$, $\overline{\mathbf{16}}_H$, and the other components in $\mathbf{16}_H$, $\overline{\mathbf{16}}_H$ could be made heavy by pairing them with proper brane fields. From Eq. (47), it is straightforward to show that the SUSY vacuum corresponds to $\langle S\rangle=0$, while N_H and $\overline{N}_H$ develop VEVs of order M. [After SUSY breaking in the manner of $N=1$ SUGRA, $\langle S\rangle$ acquires a VEV of order $m_{3/2}$ (gravitino mass).] They break $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$ to the MSSM gauge group, and make the massless modes in $\Sigma_{\mathbf{10}_{-4}}$ ($\equiv\mathbf{10}_B$) and $\Sigma_{\overline{\mathbf{10}}_4}$ ($\equiv\overline{\mathbf{10}}_B$) superheavy [26]. From the last term in Eq. (47), the VEV of $\overline{N}_H$ also provides masses to the right-handed Majorana neutrinos.

Because of the presence of the soft terms, $\Sigma_{1,2}$ and $\overline{\Sigma}_{1,2}$, which carry $U(1)_{PQ}$ charges, can obtain intermediate scale VEVs of order $\sqrt{m_{3/2}M_P}$. They lead to a μ term of order $m_{3/2}$ in MSSM as desired [27,28]. Of course, the presence of $U(1)_{PQ}$ also resolves the strong CP problem [29].

The Higgs fields h_1 and $\overline{h}_1$ contain color triplets as well as weak doublets. Since the triplets in h_1 and $\overline{h}_1$ are just superheavy KK modes, a small coefficient ($\mu\sim$TeV) accompanying $h_1\overline{h}_1$ more than adequately suppresses dimension 5 operators that induce proton decay. Proton decay can still proceed via superheavy gauge bosons with masses $\approx\pi/y_c$ and are adequately suppressed [$\tau_p\sim 10^{34-36}$ yr].

IV. INFLATION AND LEPTOGENESIS

The first two terms in the superpotential (47) are ideally suited for realizing an inflationary scenario along the lines described in Refs. [1,4,6]. We will not provide any details except to note that the breaking of $U(1)_X$ takes place near the end of inflation which can lead to the appearance of cosmic strings. Since the symmetry breaking scale of $U(1)_X$ is determined from inflation to be close to 10^{16} GeV [4], the cosmic strings are superheavy and therefore not so desirable (because of potential conflict with the recent $\delta T/T$ measurements). They can be simply avoided by following the strategy outlined in Ref. [6], in which suitable nonrenormalizable terms are added to Eq. (47), such that $U(1)_X$ is broken along the inflationary trajectory and the strings are inflated away. Remarkably, the salient features of the inflationary scenario are not affected by the addition of such higher order terms. For κ somewhat smaller (larger) than 10^{-3}, n varies between 0.98 and 0.99.

The inflationary epoch ends with the decay of the oscillating fields S, N_H, $\overline{N}_H$ into the MSSM singlet (right-handed) neutrinos, whose subsequent out of equilibrium decays yield the observed baryon asymmetry via leptogenesis [8,9]. The production of the right-handed neutrinos and sneutrinos proceeds via the superpotential couplings $\nu_i^c\nu_j^c\overline{N}_H\overline{N}_H$ and $SN_H\overline{N}_H$ on B1, where $\nu_{i,j}^c$ $(i,j=1,2,3)$ denote the $SU(5)$ singlet (right-handed neutrinos) carrying nonzero $U(1)_X$ charges. Taking account of the atmospheric and solar neutrino oscillation data [11], and assuming a hierarchical pattern of neutrino masses (both left- and right-handed ones), it turns out that the inflaton decays into the lightest (the first family right handed neutrino) [1,30]. The discussion proceeds along the lines discussed in Ref. [31], where it was shown that a baryon asymmetry of the desired magnitude is readily obtained.

The 5D inflationary solution requires positive vacuum energies on both branes B1 and B2 [1]. While inflation could be driven by the first two terms in Eq. (47) at B1, an appropriate scalar potential on B2 is also necessary. Since the boundary conditions in Eqs. (11) and (12) require Λ_1 and Λ_2 to simultaneously vanish, it is natural for S to be a bulk field. The VEVs of S on the two branes can be adjusted such that the boundary conditions are satisfied. As a simple example, consider the following superpotential at B2,

$$W_{B2}=\kappa_2 S(Z\overline{Z}-M_2^2), \tag{48}$$

where Z and $\overline{Z}$ are $SO(10)$ singlet superfields with opposite $U(1)_R$ charges. Thus, only the gauge symmetry $U(1)_X$ on B1 is broken during inflation.

V. CONCLUSION

Inspired by recent attempts to construct realistic 5D SUSY GUT models, we have presented a realistic inflationary scenario in this setting, along the lines proposed in Ref. [1]. Inflation is implemented through F-term scalar potentials on the two branes, which is allowed by 5D Einstein gravity. We have discussed the transition from the inflationary to the radiation dominated phase, and provided a realistic 5D SUSY $SO(10)$ model in which the compactification on an $S^1/(Z_2\times Z_2')$ orbifold leads to the gauge symmetry $SU(3)_c\times SU(2)_L\times U(1)_Y\times U(1)_X$. Inflation is associated with the breaking of the gauge symmetry $U(1)_X$, at a scale very close to $M_{\rm GUT}$. Baryogenesis via leptogenesis is very natural in this approach, and the scalar spectral index $n=0.98$–0.99. The gravitational contribution to the quadrupole anisotropy is found to be tiny ($\lesssim 10^{-4}$).

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