Leptogenesis in Smooth Hybrid Inflation

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\textbf{Abstract}

We present a concrete supersymmetric grand unified model based on the Pati-Salam gauge group $SU(4)_c \times SU(2)_L \times SU(2)_R$ and leading naturally to smooth hybrid inflation, which avoids the cosmological disaster encountered in the standard hybrid inflationary scenario from the overproduction of monopoles at the end of inflation. Successful ‘reheating’ which satisfies the gravitino constraint takes place after the termination of inflation. Also, adequate baryogenesis via a primordial leptogenesis occurs consistently with the solar and atmospheric neutrino oscillation data as well as the $SU(4)_c$ symmetry.

The most promising theory for solving the cosmological magnetic monopole problem of grand unified theories (GUTs) is inflation. This theory also solves many of the outstanding problems of the standard Big Bang cosmological model and predicts the formation of the large scale structure in the universe as well as the temperature fluctuations which are observed in the cosmic microwave background radiation (CMBR). However, the early realizations of inflation require extremely flat potentials and very small coupling constants. To solve this naturalness problem, Linde has introduced \cite{1}, some years ago, the hybrid inflationary scenario which involves two real scalar fields, one of which may be a gauge non-singlet. Unfortunately, in this scheme, the (abrupt) termination of inflation is followed by a ‘waterfall’ regime during which topological defects can be copiously produced \cite{2}.
In particular, the cosmological problem caused by the overproduction of GUT magnetic monopoles is not avoided in this inflationary scenario.

The simplest framework for realizing hybrid inflation is provided \cite{3,4} by the supersymmetric (SUSY) GUTs which are based on gauge groups with rank greater than five. The same superpotential which breaks spontaneously the GUT gauge group to a subgroup with lower rank also leads \cite{4} to successful hybrid inflation with ‘natural’ values of the relevant coupling constant. The slowly rolling inflaton field belongs to a gauge singlet superfield which couples to a conjugate pair of gauge non-singlet Higgs superfields. These fields acquire non-vanishing vacuum expectation values (vevs) after the end of inflation, thereby breaking the GUT gauge symmetry. The tree-level scalar potential possesses a flat valley of local minima for values of the gauge singlet inflaton greater than a certain critical value. The vevs of the Higgs superfields vanish along this valley. Inflation takes place when the system is trapped in the valley of local minima. The (classical) flatness of the valley is lifted by the one-loop radiative corrections \cite{4} to the scalar potential. They are non-zero because of the SUSY breaking caused by the non-vanishing vacuum energy density on the valley. The slow-roll conditions (see e.g., Ref.\cite{5}) are satisfied for all values of the gauge singlet inflaton greater than its critical value, where inflation ends abruptly. It is followed by a ‘waterfall’ regime during which the Higgs fields quickly fall towards the SUSY minima of the potential and oscillate about them. If the SUSY vacuum manifold is homotopically non-trivial, topological defects will be copiously formed by the Kibble mechanism since the system can end up at any point of the vacuum manifold with equal probability. So a cosmological disaster is encountered in the hybrid inflationary models which are based on GUTs predicting the existence of magnetic monopoles.

The monopole problem can be solved by utilizing the following observation \cite{6}. The standard superpotential of SUSY hybrid inflation involves only renormalizable terms. Note, though, that an infinite number of non-renormalizable terms, which are linear in the gauge singlet inflaton, cannot be excluded by any symmetries. These terms are usually neglected. However, the leading of these terms, if its dimensionless coefficient is of order unity, can compete with the trilinear term of the standard superpotential whose coupling constant is typically of order $10^{-3}$. A ‘shifted’ classically flat valley of local minima where the GUT gauge symmetry is broken then also appears \cite{6} and can be used as an alternative inflationary trajectory. The necessary inclination along this valley is again given by the one-loop radiative corrections to the potential which are now calculated \cite{6} with both the
gauge singlet inflaton and the Higgs superfields acquiring constant non-zero values. This
scheme is known [7] as shifted hybrid inflationary scenario. Inflation again ends abruptly
by a ‘waterfall’ with the system falling into the SUSY vacua. The main difference from
standard (SUSY) hybrid inflation is that now no topological defects can form at the end
of inflation since the GUT gauge symmetry is already broken during inflation.

In Ref.[6], a concrete SUSY GUT model based on the Pati-Salam (PS) gauge group
\[ G_{PS} = SU(4)_c \times SU(2)_L \times SU(2)_R \]
leading to successful shifted hybrid inflation has been constructed. The SUSY PS model is particularly interesting since it can also
be obtained [8] from string compactifications. It is one of the simplest GUT models
predicting the existence of magnetic monopoles. These PS monopoles, which carry two
units of ‘Dirac’ magnetic charge [9], would then lead to a cosmological disaster if standard
hybrid inflation was employed. The results of the cosmic background explorer (COBE) [10]
were reproduced in this model with ‘natural’ values of the relevant parameters and \( G_{PS} \)
spontaneous breaking scale of the order of \( 10^{16} \) GeV. Moreover, the \( \mu \) term was generated
via a Peccei-Quinn (PQ) symmetry and proton was practically stable. Hierarchical light
neutrino masses were generated via the seesaw mechanism. Baryogenesis in the universe
could occur via a primordial leptogenesis [11] consistently with the gravitino limit [12] on
the ‘reheat’ temperature, the \( SU(4)_c \) symmetry and the experimental data on solar and
atmospheric neutrino oscillations.

An alternative solution to the monopole problem of (SUSY) hybrid inflation had been
given some years ago in Ref.[2]. The idea was to impose an extra discrete \( Z_2 \) symmetry
which forbids the (renormalizable) trilinear coupling of the standard inflationary super-
potential and introduce instead the leading non-renormalizable term. The inflationary
superpotential then coincides with the ‘shifted’ one, but without the trilinear term. The
emerging picture is, however, dramatically different from the one encountered in shifted
hybrid inflation. The scalar potential possesses two symmetric valleys of local minima
which are suitable for inflation and along which the GUT gauge symmetry is broken. The
inclination of these valleys, which contain the SUSY minima too, is already non-zero at
the classical level and thus the one-loop radiative corrections to the scalar potential are
not needed. In Ref.[13], it has been shown that, for almost all initial conditions, the system
is trapped in either of these valleys and generates an adequate number of e-foldings
as it moves towards the SUSY minimum following [2] the bottom of the valley. Inflation
ends when the slow-roll conditions are violated and the system smoothly enters into an
oscillatory phase about the SUSY minimum. This is the reason for calling this scenario smooth hybrid inflation. Since a specific inflationary path leading to a specific SUSY minimum is chosen already from the beginning of inflation, no topological defects can form at the end of inflation. In particular, the cosmological monopole problem is solved. The CMBR quadrupole anisotropy can be reproduced with ‘natural’ values of the parameters and a gauge symmetry breaking scale identical to the SUSY GUT scale.

In this letter, we present a concrete SUSY GUT model which naturally leads to smooth hybrid inflation. It is based on the PS gauge group and is consistent with all the cosmological and phenomenological requirements. We pay particular attention to the study of the ‘reheating’ process following inflation and the generation of the observed baryon asymmetry of the universe (BAU) via a primordial leptogenesis [11] in this model.

We start with the PS SUSY GUT model of Ref.[6]. The breaking of $G_{PS}$ down to the standard model gauge group is achieved by the superheavy vevs of the right-handed neutrino components of a conjugate pair of Higgs superfields

$$
\tilde{H}^c = (4, 1, 2) \equiv \left( \begin{array}{c} \tilde{u}_H^c \\ \tilde{d}_H^c \\ \tilde{c}_H^c \\ \tilde{\nu}_H^c \end{array} \right),
$$

$$
H^c = (\bar{4}, 1, 2) \equiv \left( \begin{array}{c} u_H^c \\ d_H^c \\ c_H^c \end{array} \right).
$$

The three families of ‘matter’ superfields are $F_i = (4, 2, 1)$ and $F^c_i = (\bar{4}, 1, 2) \ (i = 1, 2, 3)$, while the electroweak Higgs doublets $h^{(1)}, h^{(2)}$ belong to the superfield $h = (1, 2, 2)$. The model also contains a gauge singlet $S$ which triggers the breaking of $G_{PS}$, an $SU(4)_c$ 6-plet $G = (6, 1, 1)$ which gives masses to $\tilde{d}_H^c, \tilde{d}_H^c$ (see first paper in Ref.[8]), and a pair of gauge singlets $\tilde{N}, N$ for solving [14] the $\mu$ problem of the minimal SUSY standard model (MSSM) via a PQ symmetry which also solves the strong CP problem. In addition to $G_{PS}$, the model possesses two global $U(1)$ symmetries, namely a PQ and a R-symmetry, as well as a discrete $Z^{mp}_2$ symmetry (‘matter parity’) under which $F, F^c$ change sign. For details on the charge assignments and the full superpotential, the reader is referred to Ref.[6].

Baryon and lepton number are not conserved in this model and proton can decay via effective dimension five operators from one-loop graphs. However, its lifetime turns out to be long enough so that it is practically stable. The right-handed neutrinos $\nu^c_i$ acquire intermediate Majorana masses via non-renormalizable couplings to the superheavy vevs of $\tilde{H}^c, H^c$. Light neutrino masses are then generated via the seesaw mechanism.
We impose an extra $Z_2$ symmetry under which $H^c \rightarrow -H^c$. The structure of the model remains unaltered except that now only even powers of the combination $\tilde{H}^c H^c$ are allowed in the superpotential terms. In particular, all the properties summarized in the previous paragraph are still valid. The only physically significant alteration is that the trilinear term $S\tilde{H}^c H^c$ is now missing from the inflationary superpotential. Including instead the leading non-renormalizable term, this superpotential becomes

$$\delta W = S \left(-\mu^2 + \frac{(\tilde{H}^c H^c)^2}{M_S^2}\right).$$

(2)

Here $\mu$ is a superheavy mass parameter and $M_S \sim 5 \times 10^{17}$ GeV is the string mass scale. The dimensionless coupling constants have been absorbed in $\mu$ and $M_S$. Note that these two mass parameters can be made positive by field redefinitions.

The inflationary scalar potential $V$ derived from $\delta W$ in Eq.(2) is given by

$$\tilde{V} = \frac{V}{\mu^4} = (1 - \tilde{\chi}^4)^2 + 16\tilde{\sigma}^2\tilde{\chi}^6,$$

(3)

where we have used the dimensionless fields $\tilde{\chi} = \chi/2(\mu M_S)^{1/2}$ and $\tilde{\sigma} = \sigma/2(\mu M_S)^{1/2}$ with $\chi$, $\sigma$ being normalized real scalar fields defined by $\nu^c = \nu^c = \chi/2$, $S = \sigma/\sqrt{2}$ after rotating $\tilde{\nu}^c_H$, $\tilde{\nu}^c_H$, $S$ to the real axis by appropriate gauge and R-transformations. This potential has a completely different structure from both the standard and ‘shifted’ hybrid inflationary potentials. It still possesses a flat direction at $\tilde{\chi} = 0$, but this is now a local maximum with respect to $\tilde{\chi}$ for all values of $\tilde{\sigma}$. It also has two symmetric valleys of local minima with respect to $\tilde{\chi}$ at

$$\tilde{\chi} = \pm \sqrt{6\tilde{\sigma}} \left[ \left(1 + \frac{1}{36\tilde{\sigma}^4}\right)^{1/2} - 1 \right]^{1/2},$$

(4)

which can be used as inflationary trajectories. They contain the SUSY vacua which lie at $\tilde{\chi} = \pm 1$, $\tilde{\sigma} = 0$. Note that these valleys are not classically flat. In fact, already at the tree level, they possess an inclination, which can drive the inflaton towards the SUSY vacua. As a consequence, contrary to the case of standard SUSY or shifted hybrid inflation, there is no need of radiative corrections, which are expected to give a subdominant contribution to the slope of the inflationary paths. In spite of this, one could try to include the one-loop corrections. This requires the construction of the mass spectrum on the inflationary trajectories. In doing so, we find that the mass$^2$ of some scalars belonging to the inflaton sector is negative. The one-loop corrections, which involve logarithms of the masses
squared, are then ill-defined. This may be remedied by resumming the perturbative expansion to all orders, which is a formidable task and we do not pursue it here.

The dimensionless potential along the inflationary valleys is given by

$$\tilde{V} = 48\tilde{\sigma}^4 \left[ 72\tilde{\sigma}^4 \left( 1 + \frac{1}{36\tilde{\sigma}^4} \right) \left( 1 + \frac{1}{36\tilde{\sigma}^4} \right)^{\frac{1}{2}} - 1 \right] - 1 \right].$$

The system follows \[2\], from the beginning, a particular inflationary path and, thus, ends up at a specific point of the vacuum manifold leading to no production of monopoles. Also, no disastrous domain walls from the spontaneous breaking of the extra $Z_2$ symmetry under which $H^c$ changes sign are generated. The termination of inflation is not abrupt as in the other two hybrid inflationary scenarios. The reason is that the inflationary path is stable with respect to $\tilde{\chi}$ for all values of $\tilde{\sigma}$. Inflation ends smoothly at a value $\tilde{\sigma}_0$ of $\tilde{\sigma}$ after which the slow-roll conditions cease to hold. The leading slow-roll parameters $\epsilon$ and $\eta$ (see e.g., Ref.\[3\]) involve the derivatives of the potential along the inflationary path, which are

$$\frac{dV}{d\tilde{\sigma}} = 192\tilde{\sigma}^3 \left( 1 + 144\tilde{\sigma}^4 \right) \left( 1 + \frac{1}{36\tilde{\sigma}^4} \right)^{\frac{1}{2}} - 1 \right] - 2 \right],$$

$$\frac{d^2V}{d\tilde{\sigma}^2} = \frac{16}{3\tilde{\sigma}^2} \left( 1 + 504\tilde{\sigma}^4 \right) \left[ 72\tilde{\sigma}^4 \left( 1 + \frac{1}{36\tilde{\sigma}^4} \right)^{\frac{1}{2}} - 1 \right] - 1 \right]$$

$$- \left( 1 + 252\tilde{\sigma}^4 \right) \left( 1 + \frac{1}{36\tilde{\sigma}^4} \right)^{-\frac{1}{2}} - 1 \right] \right].$$

We calculate the quadrupole anisotropy of the CMBR and the number of e-foldings, $N_Q$, of our present horizon during inflation using the standard formulae (see e.g., Ref.\[3\]) and Eq.(6). The present inflationary scenario has an important advantage. The common vev of $\bar{H}^c$ and $H^c$ at the SUSY vacuum, which equals $(\mu M_S)^{1/2}$, is not so tightly restricted as in the previous hybrid inflationary scenarios. Using this freedom, we choose it equal to the SUSY GUT scale $M_G \approx 2.86 \times 10^{16}$ GeV. From the results of COBE [10] and for $N_Q \approx 57$, we then obtain $M_S \approx 4.39 \times 10^{17}$ GeV and $\mu \approx 1.86 \times 10^{15}$ GeV, which are quite ‘natural’. The value of $\sigma$ at which inflation ends corresponds to $\eta = -1$ and is $\sigma_0 \approx 1.34 \times 10^{17}$ GeV ($\epsilon$ remains always much smaller than unity). Finally, the value of $\sigma$ at which our present horizon crosses outside the inflationary horizon is $\sigma_Q \approx 2.71 \times 10^{17}$ GeV.

We now turn to the discussion of the ‘reheating’ process following inflation and the generation of the observed BAU. After the end of inflation, the system smoothly enters
into a phase of damped oscillations about the SUSY vacuum. The oscillating inflaton fields are two complex scalars \( \theta = (\delta \tilde{\nu}^c_H + \delta \nu^c_H) / \sqrt{2} \) (\( \delta \tilde{\nu}^c_H, \delta \nu^c_H \) are the deviations of \( \tilde{\nu}^c_H, \nu^c_H \) from their common vev) and \( S \). They have the same mass \( m_{\text{infl}} = \frac{2}{\sqrt{2}} \mu^2 / M_G \approx 3.42 \times 10^{14} \text{ GeV} \) and decay into right-handed neutrinos and sneutrinos respectively via the superpotential terms \( \gamma_i \tilde{H}^c \tilde{F}_i^c \tilde{F}_i^c / M_S \) (in a basis where the \( \gamma \)'s are diagonal) and \( S(\tilde{H}^c \tilde{H}^c)^2 / M_S^2 \). Their common decay width is

\[
\Gamma = \frac{1}{8\pi} \left( \frac{M_i}{M_G} \right)^2 m_{\text{infl}},
\]

where \( M_i = 2\gamma_i \mu \) is the mass of the heaviest \( \nu^c_i \) satisfying the inequality \( M_i < m_{\text{infl}} / 2 \). To minimize the number of small coupling constants, we assume that

\[
M_2 < \frac{1}{2} m_{\text{infl}} \leq M_3 = 2\gamma_3 \mu,
\]

so that the oscillating inflaton fields decay into the second heaviest right-handed neutrino superfield \( \nu^c_2 \). For MSSM spectrum, the ‘reheat’ temperature \( T_r \approx (1/7)(\Gamma M_P)^{1/2} = (1/7)(M_2/M_G)(m_{\text{infl}} M_P / 8\pi)^{1/2} \) [15], where \( M_P \approx 1.22 \times 10^{19} \text{ GeV} \) is the Planck mass scale. Hence, \( M_2 \) can be expressed in terms of \( T_r \). The subsequent decay of \( \nu^c_2 \) into lepton and electroweak Higgs superfields creates a primordial lepton asymmetry [11] which survives [16] until the electroweak transition, where it is converted in part into baryon asymmetry via non-perturbative electroweak sphaleron effects.

Analysis [17] of the CHOOZ experiment [18] shows that the solar and atmospheric neutrino oscillations decouple, allowing us to concentrate on the two heaviest families. The light neutrino mass matrix, which is generated via the seesaw mechanism, is

\[
m_\nu \approx -\tilde{m}^D \frac{1}{M^R} m^D,
\]

where \( m^D \) is the ‘Dirac’ neutrino mass matrix and \( M^R \) the Majorana mass matrix of right-handed neutrinos with (positive) eigenvalues \( m^D_{2,3} \) and \( M_{2,3} \) respectively. The two (positive) eigenvalues of \( m_\nu \) are denoted by \( m_2 = m_\nu_\mu \) and \( m_3 = m_\nu_\tau \). The determinant and trace invariance of \( m_\nu^T m_\nu \) provide us with two constraints [19] on the mass parameters \( m_{2,3}, m^D_{2,3}, M_{2,3} \) and the rotation angle \( \theta \) and phase \( \delta \) which diagonalize \( M^R \) in the basis where \( m^D \) is diagonal. The primordial lepton asymmetry is given [19] by

\[
\frac{n_L}{s} \approx 1.33 \frac{9T_r}{16\pi m_{\text{infl}}} \frac{M_2}{M_3} \frac{c^2 s^2 \sin 2\delta (m_3^D - m_2^D)^2}{|\langle h^{(1)} \rangle|^2 (m_3^D s^2 + m_2^D c^2)^2},
\]

where \( c = \cos \theta, s = \sin \theta \) and the electroweak vev coupling to the up-type quarks \( |\langle h^{(1)} \rangle| \approx 174 \text{ GeV} \) since, here, \( \tan \beta \) is large as implied [20] by the fact that the electroweak Higgs as
well as the right-handed quark superfields form $SU(2)_R$ doublets. For MSSM spectrum, $n_L/s \approx -(79/28)(n_B/s)$, where $n_B/s$ is the BAU (see second paper in Ref.[10]). Note that Eq.(11) holds provided [21] that $M_2 \ll M_3$ and the decay width of $\nu_3^c$ is $\ll (M_3^2 - M_2^2)/M_2$, which are well-satisfied here. The $\mu - \tau$ mixing angle $\theta_{23} = \theta_{\mu\tau}$ lies [19] in the range

$$|\varphi - \theta^D| \leq \theta_{\mu\tau} \leq \varphi + \theta^D, \quad \text{for} \quad \varphi + \theta^D \leq \pi/2,$$

(12)

where $\varphi$ is the rotation angle which diagonalizes $m_\nu$ in the basis where $m^D$ is diagonal and $\theta^D$ is the ‘Dirac’ mixing angle defined with vanishing Majorana masses for the $\nu^c$’s.

We are now ready to examine whether the gravitino constraint [12] on the ‘reheat’ temperature and the restrictions on the BAU from Big Bang nucleosynthesis can be satisfied consistently with the data on solar and atmospheric neutrino oscillations and the $SU(4)_c$ symmetry. The gravitino constraint is usually quoted as $T_r \lesssim 10^9 \text{ GeV}$. Unfortunately, we do not find any solutions for such values of $T_r$. Therefore, we take $T_r = 10^{10} \text{ GeV}$, which is also perfectly acceptable provided [22] that the branching ratio of the gravitino to photons is somewhat smaller than unity and the gravitino mass is relatively large (of order a few hundred GeV). The low deuterium abundance constraint [23] on the BAU, $0.017 \lesssim \Omega_B h^2 \lesssim 0.021$, yields the bound $1.8 \times 10^{-10} \leq -n_L/s \lesssim 2.3 \times 10^{-10}$.

The small or large mixing angle MSW solution of the solar neutrino puzzle requires $2 \times 10^{-3} \text{ eV} \leq m_{\nu_2} \lesssim 3.2 \times 10^{-3} \text{ eV}$ or $3.6 \times 10^{-3} \text{ eV} \leq m_{\nu_2} \lesssim 13 \times 10^{-3} \text{ eV}$ respectively [24]. The $\tau$-neutrino mass is restricted in the range $3 \times 10^{-2} \text{ eV} \leq m_{\nu_2} \lesssim 9 \times 10^{-2} \text{ eV}$ from the results of the SuperKamiokande experiment [27] which also imply almost maximal $\nu_\mu - \nu_\tau$ mixing, i.e., $\sin^2 2\theta_{\mu\tau} \gtrsim 0.85$. We assume, for simplicity, that the ‘Dirac’ mixing angle $\theta^D$ is negligible, so that $\theta_{\mu\tau} \approx \varphi$. The $SU(4)_c$ symmetry implies that $m_3^D$ coincides asymptotically with the top quark mass. Taking renormalization effects into account, for MSSM spectrum with large $\tan \beta$, we obtain [19] $m_3^D \approx 110 - 120 \text{ GeV}$. We also include the running of $\theta_{\mu\tau}$ from $M_G$ to the electroweak scale (see last paper in Ref.[11]).

Our results are shown in Figs.1, 2 and 3 where we plot solutions corresponding to $T_r = 10^{10} \text{ GeV}$ and satisfying the leptogenesis constraint consistently with the neutrino oscillation data and the $SU(4)_c$ symmetry. The parameter $\gamma_3$ runs from 0.05 to 0.5, i.e., $M_3 \approx 1.86 \times 10^{14} - 1.86 \times 10^{15} \text{ GeV}$. The second inequality in Eq.(3) implies that $\gamma_3 \gtrsim 0.046$. However, no solutions are found for $\gamma_3 < 0.05$. Also, values of $\gamma_3$ higher than 0.5 do not allow solutions. The mass of the second heaviest right-handed neutrino $M_2 \approx 1.55 \times 10^{11} \text{ GeV}$, which corresponds to $T_r = 10^{10} \text{ GeV}$ and clearly satisfies the
first inequality in Eq. (9). The restrictions from $SU(4)_c$ invariance are expected to be more or less accurate only if applied to the masses of the third family quarks and leptons. For the second family, they should hold only as order of magnitude relations. We thus restrict ourselves to values of $m_2^D$ smaller than 2 GeV since much bigger $m_2^D$’s would violate strongly the $SU(4)_c$ symmetry (the value of $m_2^D$ from exact $SU(4)_c$ is about 0.23 GeV for MSSM spectrum with large tan $\beta$). Moreover, we find that solutions exist only if $m_2^D \gtrsim 0.8$. So we take $m_2^D \approx 0.8 - 2$ GeV and, as required by $SU(4)_c$ invariance, $m_3^D \approx 110 - 120$ GeV. Also, the phase $\delta \approx (-\pi/8) - (-\pi/5)$ and the rotation angle $\theta \approx 0.01 - 0.03$ for solutions to appear. Note that $\delta$’s close to 0 or $-\pi/2$ are excluded since they yield very small primordial lepton asymmetry.

In Fig. 1, we present the scatter plot of our solutions in the $m_{\nu_\tau} - \sin^2 2\theta_{\mu\tau}$ plane. In Figs. 2 and 3, solutions are plotted in the $m_{\nu_\mu} - \sin^2 2\theta_{\mu\tau}$ and $m_{\nu_\tau} - m_{\nu_\mu}$ planes respectively. Our model favors the large mixing angle MSW solution to the solar neutrino problem. For large $m_{\nu_\tau}$’s (close to the upper bound) or small $m_{\nu_\mu}$’s (close to the lower bound), it excludes very large $\mu - \tau$ mixing and, in any case, disfavors large $m_{\nu_\tau}$’s.

In conclusion, we presented a concrete PS SUSY GUT model leading naturally to smooth hybrid inflation, which solves the cosmological magnetic monopole problem. The model reproduces the measured quadrupole anisotropy of the CMBR with ‘natural’ values of the parameters and a PS spontaneous breaking scale equal to the SUSY GUT scale. A PQ symmetry is used to generate the $\mu$ term of MSSM and proton is practically stable. Inflation is followed by a successful ‘reheating’ process satisfying the gravitino constraint on the ‘reheat’ temperature and generating the observed BAU via a primordial leptogenesis consistently with the requirements from solar and atmospheric neutrino oscillations and the $SU(4)_c$ symmetry.

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Figure 1: The scatter plot in the $m_{\nu_\tau} - \sin^2 2\theta_{\mu\tau}$ plane of the solutions which satisfy the low deuterium abundance constraint on the BAU, the restrictions from solar and atmospheric neutrino oscillations, and the $SU(4)_c$ invariance. We take $T_r = 10^{10}$ GeV, $\gamma_3 \approx 0.05 - 0.5$, $m_2^D \approx 0.8 - 2$ GeV and $m_3^D \approx 110 - 120$ GeV.
Figure 2: The scatter plot in the $m_{\nu_{\mu}} - \sin^2 2\theta_{\mu\tau}$ plane of the solutions depicted in Fig.[1].

Figure 3: The scatter plot in the $m_{\nu_{\tau}} - m_{\nu_{\mu}}$ plane of the solutions depicted in Fig.[4].