Atmospheric Neutrino Anomaly and Supersymmetric Inflation

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Abstract

A detailed investigation of hybrid inflation and the subsequent reheating process is performed within a $\mu$-problem solving supersymmetric model based on a left-right symmetric gauge group. The process of baryogenesis via leptogenesis is especially studied. For $\nu_\mu$, $\nu_\tau$ masses from the small angle MSW resolution of the solar neutrino problem and the recent results of the SuperKamiokande experiment, we show that maximal $\nu_\mu$-$\nu_\tau$ mixing can be achieved. The required value of the relevant coupling constant is, however, quite small ($\sim 10^{-6}$).

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The hybrid inflationary scenario \cite{1} can be easily implemented \cite{2-4} in the context of supersymmetric theories in a ‘natural’ way meaning that a) there is no need for tiny coupling constants, b) the superpotential used is the most general one allowed by gauge and R- symmetries, c) supersymmetry guarantees that radiative corrections do not invalidate inflation, but rather provide a slope along the inflationary trajectory which drives the inflaton towards the supersymmetric vacua, and d) supergravity corrections can be brought under control so as to leave inflation intact.

A moderate extension of the minimal supersymmetric standard model (MSSM) based on a left-right symmetric gauge group provides \cite{4} a suitable framework for hybrid inflation. The inflaton is associated with the breaking of $SU(2)_R$ and consists of a gauge singlet and a pair of $SU(2)_R$ doublets. The doublets can decay into right handed neutrinos, after inflation, reheating the universe and providing a mechanism \cite{5} for baryogenesis through a primordial leptogenesis. The gauge singlet, however, has no direct coupling to light matter in the simplest case. Moreover, its coupling to the $SU(2)_R$ doublets turns out to be unable to ensure its efficient decay. This difficulty can be overcome by introducing \cite{4,6} a direct superpotential coupling of the gauge singlet superfield to the electroweak higgs doublets. This way the gauge singlet scalar can decay into a pair of higgsinos. It has been shown \cite{6} that, in the presence of gravity-mediated supersymmetry breaking, this gauge singlet acquires a vacuum expectation value (vev) and consequently generates, through its coupling to the ordinary higgs superfields, the $\mu$ term of MSSM.

A coupling of the scalar components of the $SU(2)_R$ doublets to the electroweak higgses is automatically induced in this scheme, allowing them to decay into a pair of ordinary higgses in addition to their useful (for baryogenesis) decay to right handed neutrinos.

In this paper, we attempt a detailed study of inflation in the above scheme. In particular, we solve the evolution equations of this system and estimate the reheating temperature. The process of baryogenesis via leptogenesis is also considered and its consequences on $\nu_\mu$-$\nu_\tau$ mixing are analyzed. For masses of $\nu_\mu$, $\nu_\tau$ which are consistent with the small angle MSW resolution of the solar neutrino problem and the recent results of the SuperKamiokande experiment \cite{7}, we examine whether maximal $\nu_\mu$-$\nu_\tau$ mixing can be achieved.

Let us first describe the main features of the $G_{LR} = SU(3)_c \times SU(2)_R \times SU(2)_L \times$
U(1)$_{B-L}$ symmetric model which solves the $\mu$ problem. The $SU(2)_R \times U(1)_{B-L}$ group is broken by a pair of $SU(2)_R$ doublet chiral superfields $l^c$, $\bar{l}^c$ which acquire a vev $M \gg m_{3/2} \sim (0.1 - 1)$ TeV, the gravitino mass. This breaking is achieved by means of a gauge singlet chiral superfield $S$ which plays a crucial three-fold role: 1) it triggers $SU(2)_R$ breaking; 2) it generates the $\mu$ term of MSSM after gravity-mediated supersymmetry breaking; and 3) it leads to hybrid inflation. Ignoring the matter fields of the model, the superpotential reads

$$W = S(\kappa l^c \bar{l}^c + \lambda h^2 - \kappa M^2),$$

where the chiral superfield $h = (h^{(1)}, h^{(2)})$ belongs to a bidoublet $(2, 2)_0$ representation of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, and $h^2$ denotes the unique bilinear invariant $\epsilon^{ij} h^{(1)}_i h^{(2)}_j$. Note that the parameters $\kappa$, $\lambda$ and $M$ can be made positive through a suitable redefinition of the superfields. $W$ in Eq.(1) has the most general renormalizable form invariant under the gauge group and a continuous $U(1)$ R-symmetry under which $S$ carries the same charge as $W$, while $h$, $l^c$, $\bar{l}^c$ are neutral. This symmetry is extended to include the matter fields of the model too and implies automatic baryon number conservation. It has been shown that, after supersymmetry breaking, $S$ develops a vev $\langle S \rangle \approx -m_{3/2}/\kappa$ which generates a $\mu$ term with $\mu = \lambda \langle S \rangle \approx -\lambda/\kappa m_{3/2}$.

The model has a built-in inflationary trajectory in the field space along which the $F_S$ term is constant. This trajectory is parametrized by $|S|$, $|S| > S_c = M$ for $\lambda > \kappa$ (see below). All other fields vanish on this trajectory. The $F_S$ term provides us with a constant tree level vacuum energy density $V_{\text{tree}} = \kappa^2 M^4$, which is responsible for inflation. Radiative corrections generate a logarithmic slope along the inflationary trajectory that drives the inflaton toward the minimum. The one-loop contribution to this slope comes from the $l^c$, $\bar{l}^c$ and $h$ supermultiplets, which receive at tree level a non-supersymmetric contribution to the masses of their scalar components from the $F_S$ term. For $|S| \leq S_c = M$, the $l^c$, $\bar{l}^c$ components become tachyonic, compensate the $F_S$ term and the system evolves towards the ‘correct’ supersymmetric minimum at $h = 0$, $l^c = \bar{l}^c = M$. (For $\kappa > \lambda$, $h$ would have become tachyonic earlier and the system would have evolved towards the ‘wrong’ minimum at $h \neq 0$, $l^c = \bar{l}^c = 0$.) Inflation can continue at least till $|S|$ approaches the instability at $|S| = S_c$ provided that the slow roll conditions

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are violated only ‘infinitesimally’ close to it. This is true for all values of the relevant parameters considered in this work. The cosmic microwave quadrupole anisotropy can be calculated [3] by standard methods and turns out to be

$$\left(\frac{\delta T}{T}\right)_Q \approx \frac{32\pi^{5/2}}{3\sqrt{3}} \left(\frac{M}{M_P}\right)^3 \kappa^{-1} x_Q^{-1} \Lambda(x_Q)^{-1},$$  

(2)

where $M_P = 1.22 \times 10^{19}$ GeV is the Planck scale and

$$\Lambda(x) = \left(\frac{\lambda}{\kappa}\right)^3 \left[\left(\frac{\lambda}{\kappa} x^2 - 1\right) \ln \left(1 - \frac{\kappa}{\lambda} x^{-2}\right) + \left(\frac{\lambda}{\kappa} x^2 + 1\right) \ln \left(1 + \frac{\kappa}{\lambda} x^{-2}\right)\right]$$  

$$+ (x^2 - 1) \ln(1 - x^{-2}) + (x^2 + 1) \ln(1 + x^{-2}),$$  

(3)

with $x = |S|/S_c$ and $S_Q$ being the value of $|S|$ when the present horizon scale crossed outside the inflationary horizon. The number of e-foldings experienced by the universe between the time the quadrupole scale exited the horizon and the end of inflation is

$$N_Q \approx 32\pi^3 \left(\frac{M}{M_P}\right)^2 \kappa^{-2} \int_1^{x_Q} \frac{dx}{x} \Lambda(x)^{-1}.$$  

(4)

The spectral index of density perturbations turns out to be very close to unity.

After reaching the instability at $|S| = S_c$, the system undergoes [9] a short complicated evolution during which inflation continues for another e-folding or so. The energy density of the system is reduced by a factor of about 2-3 during this period. The system then rapidly settles in a regular oscillatory phase about the supersymmetric vacuum. Parametric resonance can be ignored in this case [4]. The inflaton (oscillating system) consists of the two complex scalar fields $S$ and $\theta = (\delta \phi + \delta \bar{\phi})/\sqrt{2}$, where $\delta \phi = \phi - M$, $\delta \bar{\phi} = \bar{\phi} - M$, with mass $m_{infl} = \sqrt{2}\kappa M$. Here $\phi, \bar{\phi}$ are the neutral components of the superfields $l^c, \bar{l}^c$ respectively. The scalar fields $S$ and $\theta$ predominantly decay into ordinary higgsinos and higgses respectively with a common decay width $\Gamma_h = (1/16\pi)\lambda^2 m_{infl}$, as one can easily deduce from the couplings in Eq.(1). Note, however, that $\theta$ can also decay to right handed neutrinos $\nu^c$ through the non-renormalizable superpotential term $(M_{\nu^c}/2M^2)\bar{\phi}\nu^c\nu^c$, allowed by the gauge and R- symmetries of the model [4]. Here, $M_{\nu^c}$ denotes the Majorana mass of the relevant $\nu^c$. The scalar $\theta$ decays preferably into the heaviest $\nu^c$ with $M_{\nu^c} \leq m_{infl}/2$. The decay rate is given by
\[ \Gamma_{\nu^c} \approx \frac{1}{16\pi} \kappa^2 m_{infl} \alpha^2 (1 - \alpha^2)^{1/2} , \]  \hspace{1cm} (5)

where \( 0 \leq \alpha = 2M_{\nu^c}/m_{infl} \leq 1 \). The subsequent decay of these \( \nu^c \)’s gives rise to a primordial lepton number \( \mathbb{L} \). The baryon asymmetry of the universe can then be obtained by partial conversion of this lepton asymmetry through sphaleron effects.

The energy densities \( \rho_S, \rho_\theta, \) and \( \rho_r \) of the oscillating fields \( S, \theta \), and the ‘new’ radiation produced by their decay to higgsinos, higgses and \( \nu^c \)’s are controlled by the equations:

\[ \dot{\rho}_S = -(3H + \Gamma_h)\rho_S , \quad \dot{\rho}_\theta = -(3H + \Gamma_h + \Gamma_{\nu^c})\rho_\theta , \]  \hspace{1cm} (6)

\[ \dot{\rho}_r = -4H\rho_r + \Gamma_h\rho_S + (\Gamma_h + \Gamma_{\nu^c})\rho_\theta , \]  \hspace{1cm} (7)

where

\[ H = \frac{\sqrt{8\pi}}{\sqrt{3}M_P} (\rho_S + \rho_\theta + \rho_r)^{1/2} , \]  \hspace{1cm} (8)

is the Hubble parameter and overdots denote derivatives with respect to cosmic time \( t \).

We have assumed that the potential energy density is, to a good approximation, quadratic in the fields \( S \) and \( \theta \) and, thus, the oscillating inflaton system resembles the behavior of ‘matter’. Note that the second equation in Eq.(6) can be replaced by

\[ \rho_\theta(t) = \rho_S(t)e^{-\Gamma_{\nu^c}(t-t_0)} , \]  \hspace{1cm} (9)

where \( t_0 \) is the cosmic time at the onset of the oscillatory phase. The initial values are taken to be \( \rho_S(t_0) = \rho_\theta(t_0) \approx \kappa^2 M^4/6, \rho_r(t_0) = 0 \) and, for all practical purposes, we put \( t_0 = 0 \). The ‘reheat’ temperature \( T_r \) is calculated from the equation

\[ \rho_S + \rho_\theta = \rho_r = \frac{\pi^2}{30} g_s T_r^4 , \]  \hspace{1cm} (10)

where the effective number of massless degrees of freedom is \( g_s=228.75 \) for MSSM.

The lepton number density \( n_L \) produced by the \( \nu^c \)’s satisfies the evolution equation:

\[ \dot{n}_L = -3Hn_L + 2\epsilon\Gamma_{\nu^c}n_\theta , \]  \hspace{1cm} (11)

where \( \epsilon \) is the lepton number produced per decaying right handed neutrino and the factor of 2 in the second term of the rhs comes from the fact that we get two \( \nu^c \)’s for each decaying scalar \( \theta \) particle. Eq.(11) is easily integrated out to give
where \( a(t) \) is the scale factor of the universe. The first equation in Eq. (6) gives

\[
\rho_S(t) = \rho_S(t_0) \left( \frac{a(t)}{a(t_0)} \right)^{-3} e^{-\Gamma_h(t-t_0)},
\]

Combining Eqs. (12) and (13) we get the asymptotic \((t \to \infty)\) lepton asymmetry

\[
\frac{n_L(t)}{s(t)} \sim 3 \left( \frac{15}{8} \right)^{1/4} \pi^{-1/2} g_*^{-1/4} m_{\nu_{fl}}^{-1} \frac{\Gamma_{\nu^e}}{\Gamma_h + \Gamma_{\nu^e}} \rho_r^{-3/4} \rho_S e^{\Gamma_h t},
\]

is the asymptotic entropy density. For MSSM spectrum between 100 GeV and \( M \), the observed baryon asymmetry \( n_B/s \) is related \([10]\) to \( n_L/s \) by \( n_B/s = -(28/79)(n_L/s) \).

Assuming hierarchical light neutrino masses, we take \( m_{\nu_\mu} \approx 2.6 \times 10^{-3} \text{ eV} \) which is the central value of the \( \mu \)-neutrino mass coming from the small angle MSW resolution of the solar neutrino problem \([11]\). The \( \tau \)-neutrino mass will be restricted by the atmospheric anomaly \([7]\) in the range \( 3 \times 10^{-2} \text{ eV} \leq m_{\nu_\tau} \leq 11 \times 10^{-2} \text{ eV} \). Recent analysis \([12]\) of the results of the CHOOZ experiment \([13]\) shows that the oscillations of solar and atmospheric neutrinos decouple. We thus concentrate on the two heaviest families ignoring the first one. Under these circumstances, the lepton number generated per decaying \( \nu^e \) is \([8,14]\)

\[
\epsilon = \frac{1}{8\pi} g \left( \frac{M_3}{M_2} \right) \frac{c^2 s^2 \sin 2\delta \left( m_3^D - m_2^D \right)^2}{\left| \langle h^{(1)} \rangle \right|^2 \left( m_3^{D^2} s^2 + m_2^{D^2} c^2 \right)},
\]

where \( g(r) = r \ln(1+r^{-2}) \), \( |\langle h^{(1)} \rangle| \approx 174 \text{ GeV} \), \( c = \cos \theta \), \( s = \sin \theta \), and \( \theta \) \((0 \leq \theta \leq \pi/2)\) and \( \delta \) \((-\pi/2 \leq \delta < \pi/2)\) are the rotation angle and phase which diagonalize the Majorana mass matrix of \( \nu^e \)'s, \( M^R \), with eigenvalues \( M_2, M_3 \geq 0 \) in the basis where the 'Dirac' mass matrix of the neutrinos, \( M^D \), is diagonal with eigenvalues \( m_2^D, m_3^D \geq 0 \). Note that, for the range of parameters considered here, the scalar \( \theta \) decays into the second
heaviest right handed neutrino with mass $M_2 (< M_3)$ and, thus, $M_{\nu^c}$ in Eq. (3) should be identified with $M_2$. Moreover, $M_3$ turns out to be bigger than $m_{\text{infl}}/2$ as it should. We will denote the two positive eigenvalues of the light neutrino mass matrix by $m_2 (= m_{\nu^c})$, $m_3 (= m_{\nu^c})$ with $m_2 \leq m_3$. All the quantities here (masses, rotation angles and phases) are ‘asymptotic’ (defined at the grand unification scale $M_{\text{GUT}}$). The determinant and the trace invariance of the light neutrino mass matrix imply two constraints on the (asymptotic) parameters which take the form:

$$m_2 m_3 = \frac{(m_2^D m_3^D)^2}{M_2 M_3}, \quad (17)$$

$$m_2^2 + m_3^2 = \frac{(m_2^D 2c^2 + m_3^D 2s^2)^2}{M_2^2} + \frac{(m_3^D 2c^2 + m_2^D 2s^2)^2}{M_3^2} + 2\frac{(m_3^D 2 - m_2^D 2)^2 2c^2 s^2 \cos 2\delta}{M_2 M_3}. \quad (18)$$

The $\mu$-$\tau$ mixing angle $\theta_{23} (= \theta_{\mu\tau})$ lies in the range

$$|\varphi - \theta^D| \leq \theta_{23} \leq \varphi + \theta^D, \text{ for } \varphi + \theta^D \leq \pi/2,$$

$$|\varphi - \theta^D| \leq \theta_{23} \leq \pi - \varphi - \theta^D, \text{ for } \varphi + \theta^D \geq \pi/2,$$

where $\varphi$ ($0 \leq \varphi \leq \pi/2$) is the rotation angle which diagonalizes the light neutrino mass matrix, $m = -\tilde{M}^D M^{-1} R \tilde{M}^D$, in the basis where the ‘Dirac’ mass matrix is diagonal and $\theta^D$ ($0 \leq \theta^D \leq \pi/2$) is the ‘Dirac’ mixing angle in the 2-3 leptonic sector (i.e., the ‘unphysical’ leptonic mixing angle in the absence of the Majorana masses of the $\nu^c$’s).

The ‘asymptotic’ Dirac masses of $\nu_{\mu}, \nu_{\tau}$ as well as $\theta^D$ can be related to the quark sector parameters by assuming approximate $SU(4)_c$ symmetry. We obtain the asymptotic relations: $m_2^D \approx m_c$, $m_3^D \approx m_t$, $\sin \theta^D \approx |V_{cb}|$. Renormalization effects must now be taken into account. To this end, we take MSSM spectrum and large $\tan \beta \approx m_t/m_b$. The latter follows from the fact that the MSSM higgs doublets form a $SU(2)_R$ doublet. It turns out [14] that, in this case, renormalization effects can be accounted for by simply
substituting in the above formulae the following numerical values: \( m_2^D \approx 0.23 \text{ GeV} \), \( m_3^D \approx 116 \text{ GeV} \) and \( \sin \theta_D \approx 0.03 \). Also, \( \tan^2 2\theta_{23} \) increases by about 40% from \( M_{\text{GUT}} \) to \( M_Z \).

In order to predict the \( \nu_\mu - \nu_\tau \) mixing, we take a specific MSSM framework \cite{15} where the three Yukawa couplings of the third generation unify ‘asymptotically’ and, consequently, \( \tan \beta \approx m_t/m_b \). We choose the universal scalar mass (gravitino mass) \( m_3/2 \approx 290 \text{ GeV} \) and the universal gaugino mass \( M_{1/2} \approx 470 \text{ GeV} \). These values correspond \cite{16} to \( m_t(m_t) \approx 166 \text{ GeV} \) and \( m_A \) (the tree level mass of the CP-odd scalar higgs boson) = \( M_Z \). The ‘asymptotic’ higgsino mass \( \mu \) is related \cite{17} to these parameters by

\[
\frac{\mu}{m_3/2} \approx \left( \frac{M_{1/2}}{m_3/2} \right)^{-3/7} \left( 1 - \frac{Y_t}{Y_f} \right)\]

where \( Y_t = h^2_t \approx 0.91 \) is the square of the top-quark Yukawa coupling and \( Y_f \approx 1.04 \) is the weak scale value of \( Y_t \) corresponding to ‘infinite’ value at \( M_{\text{GUT}} \). For these numerical values, we obtain \( \lambda/\kappa = |\mu|/m_3/2 \approx 3.95 \) which can be substituted in Eqs.(2)-(4). These equations can then be solved, for \( (\delta T/T)_Q \approx 6.6 \times 10^{-6} \) from COBE, \( N_Q \approx 50 \) and any value of \( x_Q > 1 \). Eliminating \( x_Q \), we obtain \( M \) as a function of \( \kappa \) depicted in Fig.1. The evolution equations (6)-(8) are solved numerically for each value of \( \kappa \). The parameter \( \alpha^2 \) in Eq.(5) is taken to be equal to 2/3. This choice maximizes the decay width, \( \Gamma_{\nu c} \), of the inflaton to \( \nu^c \)'s and, thus, the subsequently produced primordial lepton asymmetry. The reheat temperature, \( T_r \), is then calculated from Eq.(10) for each value of \( \kappa \). The result is again depicted in Fig.4.

We next evaluate the lepton asymmetry. We begin by first considering the central value of \( m_{\nu_\tau} \approx 7 \times 10^{-2} \text{ eV} \) given by the SuperKamiokande experiment \cite{7} (the \( \mu \)-neutrino mass is kept fixed to its central value \( m_{\nu_\mu} \approx 2.6 \times 10^{-3} \text{ eV} \) from the MSW resolution of the solar neutrino puzzle). The mass of the second heaviest \( \nu^c \), into which the scalar \( \theta \) decays partially, is given by \( M_2 = M_{\nu_\tau} = \alpha m_{\text{inf}}/2 \) and \( M_3 \) can be found from Eq.(17). We can use the trace condition in Eq.(18) to solve for \( \delta(\theta) \) in the interval \( 0 \leq \theta \leq \pi/2 \). The expression for \( \delta(\theta) \) is subsequently substituted in Eq.(16) for \( \epsilon \). The leptonic asymmetry as a function of the angle \( \theta \) can be found from Eq.(14). To each value of \( \kappa \) correspond two values of the angle \( \theta \) satisfying the low deuterium abundance constraint \( \Omega_B h^2 \approx 0.025 \). (These values of \( \theta \) turn out to be quite insensitive to the exact value of \( n_B/s \).) The corresponding values of the rotation angle \( \varphi \), which diagonalizes the
light neutrino mass matrix, are then found and the allowed region of the mixing angle $\theta_{\mu\tau}$ in Eq.(19) is determined. Taking into account renomalization effects and superimposing all the permitted regions, we obtain the allowed range of $\sin^2 2\theta_{\mu\tau}$ as a function of $\kappa$, shown in Fig.2. We observe that maximal mixing ($\sin^2 2\theta_{\mu\tau} \approx 1$) is achieved for $1.5 \times 10^{-6} \lesssim \kappa \lesssim 1.8 \times 10^{-6}$. Also, $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$ corresponds to $1.2 \times 10^{-6} \lesssim \kappa \lesssim 3.4 \times 10^{-6}$.

The analysis above can be repeated for all values of $m_{\nu_\tau}$ allowed by SuperKamiokande. The allowed regions in the $m_{\nu_\tau}$-$\kappa$ plane for maximal $\nu_\mu$-$\nu_\tau$ mixing (bounded by the solid lines) and $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$ (bounded by the dotted lines) are shown in Fig.3. Notice that, for $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$, $\kappa \approx (0.9 - 7.5) \times 10^{-6}$ which is rather small. (Fortunately, supersymmetry protects this coupling from radiative corrections.) The corresponding values of $M$ and $T_r$ can be read from Fig.1. We find $1.3 \times 10^{15}$ GeV $\lesssim M \lesssim 2.7 \times 10^{15}$ GeV and $10^7$ GeV $\lesssim T_r \lesssim 3.2 \times 10^8$ GeV. We observe that $M$ turns out to be somewhat smaller than the MSSM unification scale $M_{GUT}$. (It is anticipated that $G_{LR}$ is embedded in a grand unified theory.) The reheat temperature, however, satisfies the gravitino constraint ($T_r \lesssim 10^9$ GeV). It should be noted that, for the values of the parameters chosen here, the lightest supersymmetric particle (LSP) is an almost pure bino with mass $m_{LSP} \approx 0.43M_{1/2} \approx 200$ GeV [18]. Its contribution to the mass of the universe turns out to be $\Omega_{LSP}h^2 \approx 1$.

In summary, we have investigated hybrid inflation and the subsequent reheating process in the framework of a $\mu$-problem solving supersymmetric model based on a left-right symmetric gauge group. The process of baryogenesis via leptogenesis is especially considered. For masses of $\nu_\mu$, $\nu_\tau$ consistent with the small angle MSW resolution of the solar neutrino problem and the recent SuperKamiokande data, we showed that maximal $\nu_\mu$-$\nu_\tau$ mixing can be achieved. The required value of the coupling constant $\kappa$ is, however, quite small ($\sim 10^{-6}$).

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FIG. 1. The mass scale $M$ (solid line) and the reheat temperature $T_r$ (dashed line) as functions of $\kappa$. 
FIG. 2. The allowed region (bounded by the solid lines) in the $\kappa$-$\sin^2 2\theta_{\mu\tau}$ plane for $m_{\nu_{\mu}} \approx 2.6 \times 10^{-3} \text{ eV}$ and $m_{\nu_{\tau}} \approx 7 \times 10^{-2} \text{ eV}$.
FIG. 3. The regions on the $m_{\nu_r}$-$\kappa$ plane corresponding to maximal $\nu_\mu$-$\nu_\tau$ mixing (bounded by the solid lines) and $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$ (bounded by the dotted lines). Here we consider the range $3 \times 10^{-2}$ eV $\lesssim m_{\nu_r} \lesssim 11 \times 10^{-2}$ eV ($m_{\nu_\mu} \approx 2.6 \times 10^{-3}$ eV).