A Constraint on Yukawa-Coupling Unification from Lepton-Flavor Violating Processes

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Abstract

We present a new constraint on a lepton mixing matrix V from lepton-flavor violating (LFV) processes in supersymmetric standard models with massive neutrinos. Here, we assume Yukawa-coupling unification $f_{\nu 3} \simeq f_{\text{top}}$, in which τ-neutrino Yukawa coupling $f_{\nu 3}$ is unified into top-quark Yukawa coupling f_{top} at the unification scale $M_* \simeq 3 \times 10^{16}$ GeV. We show that the present experimental bound on $\mu \to e\gamma$ decay already gives a stringent limit on the lepton mixing (typically $V_{13} < 0.02$ for $V_{23} = 1/\sqrt{2}$). Therefore, many existing neutrino-mass models are strongly constrained. Future improvement of bounds on LFV processes will provide a more significant impact on the models with the Yukawa-coupling unification. We also stress that a precise measurement of a neutrino mixing $(V_{MNS})_{e3}$ in future neutrino experiments would be very important, since the observation of non-zero $(V_{MNS})_{e3}$, together with negative experimental results for the LFV processes, have a robust potential to exclude a large class of SUSY standard models with the Yukawa-coupling unification.

1   Introduction

It is a remarkable fact that three gauge coupling constants of the SU(3)×SU(2) ×U(1) theory meet at a very high energy scale $\mu \simeq 3 \times 10^{16}$ GeV in the supersymmetric (SUSY) standard model. The grand unified theory (GUT) is the most manifest candidate to explain the unification of the three gauge coupling constants. The GUT was considered as a necessary scheme to maintain the gauge-coupling unification up to the Planck scale $M_{\text{Planck}} \simeq 2 \times 10^{18}$ GeV. However, it has been pointed out by Witten some time ago [1] that the fundamental scale can be much lower than the Planck scale in the strongly coupled string (M) theory [2] and the unification scale, $\mu \simeq 3 \times 10^{16}$ GeV, is regarded as the cut-off scale M_* of the low-energy effective field theory. That is, the standard-model gauge interactions are directly unified with gravity without going through the GUT phase. This new interpretation of gauge coupling unification has various phenomenological merits; in particular, it does not suffer from the doublet–triplet splitting problem, unlike the SUSY GUT, and it may provide a natural Peccei–Quinn axion [3] to solve the strong CP problem since the effects of world-sheet instantons are expected to be suppressed in the strongly coupled string theories [4]. In this new unification approach, however, Yukawa coupling constants are free parameters and hence we need a principle to understand another success of the SUSY GUT, i.e. $m_\tau \simeq m_b$. The Yukawa-coupling unification is the most well-known principle to explain it. Thus, we assume the Yukawa-coupling unification for the third family at the unification (cut-off) scale $M_* \simeq 3 \times 10^{16}$ GeV, and consider that the Yukawa couplings for the first and second families receive easily large threshold effects from heavy particles at the cut-off scale, since their tree-level values themselves are small compared with the Yukawa coupling constants for the third family. This interesting principle may be extended if there are right-handed neutrinos. This is because we may have another "Yukawa-coupling unification $(f_{\nu_3} \simeq f_{\text{top}})$", where f_{ν_3} is the largest eigenvalue of the Dirac mass Yukawa coupling for the neutrino.

Recently, Super-Kamiokande experiments on atmospheric neutrinos [5] have presented very convincing evidence for the oscillation of ν_μ to ν_τ with a mass difference $\delta m^2 \simeq 10^{-3} - 10^{-2}$ eV^2, which implies the largest mass of the neutrino to be $m_{\nu_3} \simeq 0.03 - 0.1$ eV, provided there is a mass hierarchy $m_{\nu_3} \gg m_{\nu_2}$. If neutrinos are indeed massive, the seesaw mechanism [6] is the most natural framework to account for the smallness of neutrino masses, where the small masses are low-energy consequences of the presence of superheavy right-

handed neutrinos. The Majorana masses of the right-handed neutrinos are determined by the Dirac mass term for neutrinos, m_{ν_D}. It is very interesting that the Yukawa-coupling unification ($f_{\nu_3} \simeq f_{\rm top}$) suggests the Majorana mass M_R of the right-handed neutrino to be $M_R \simeq 2 \times 10^{14}$ GeV, which is very close to the unification (cut-off) scale M_*.

In the models with the Yukawa-coupling unification $f_{\nu_3} \simeq f_{\rm top}$, large lepton-flavor violation (LFV) in the charged lepton sector is expected in the SUSY standard model, since the τ neutrino Yukawa coupling is very large and it induces non-negligible mass splitting among sleptons through radiative corrections [7, 8]. In this letter we perform a detailed analysis of LFV processes assuming the Yukawa-coupling unification ($f_{\nu_3} \simeq f_{\rm top}$). Throughout this paper we assume that all squarks and sleptons have a common SUSY-breaking soft mass at the unification scale $M_* \simeq 3 \times 10^{16}$ GeV taking the gravity-mediated SUSY breaking model.

We show that the present experimental upper bound on $\mu \to e\gamma$ decay already gives a stringent constraint on the mixing $V_{13}V_{23}$ (typically $V_{13} < 0.02$ for $V_{23} = 1/\sqrt{2}$). Since the constraint is very severe, the models with the Yukawa-coupling unification need an explanation (e.g. symmetry) for the smallness of V_{13}. In the existing literature, this constraint has not been taken into account. Therefore, most of the models with the Yukawa-coupling unification should be subject to this new constraint. However, we also stress that the above intriguing principle, i.e. Yukawa-coupling unification, is not yet ruled out. Future improvement of the branching ratios for the LFV processes [9, 10, 11] will provide a more significant impact on the Yukawa-coupling unification models, otherwise they will indeed be observed.

Furthermore, in most of the neutrino-mass models proposed so far, the mixing V_{13} approximately equals the neutrino mixing $(V_{MNS})_{e3}$ [12, 13, 14]. Therefore, a precise measurement of $(V_{MNS})_{e3}$ in future neutrino experiments [15, 16] is very important, since the observation of non-zero $(V_{MNS})_{e3}$ together with the negative result of $\mu \to e\gamma$ decay has a great potential to exclude a large class of SUSY standard models with the Yukawa-coupling unification.

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2 SUSY standard model with right-handed neutrinos

First, we briefly review the SUSY standard model with right-handed neutrinos. Introducing the right-handed neutrinos, we have the following superpotential in the lepton sector:

$$W = \bar{E}_i f_e^i L_i H_d + \bar{N}_i f_\nu^{ij} L_j H_u + \frac{1}{2} \bar{N}_i M_{Ri} \bar{N}_i + \text{h.c.,} \tag{1}$$

where $\bar{E}_i$, L_i and $\bar{N}_i$ are right-handed charged leptons, lepton doublets and right-handed neutrinos, respectively. In this letter, without loss of generality, we take a basis where the charged-lepton Yukawa couplings (f_e) and right-handed neutrino masses, M_R, are diagonalized.

Assuming $M_{Ri} \gg m_Z$, we integrate out the right-handed neutrinos and obtain the following effective superpotential:

$$W \simeq \bar{E}_i f_e^{ij} L_j H_d - \frac{1}{2} (f_\nu^T M_R^{-1} f_\nu)^{ij} (L_i H_u)(L_j H_u) + \text{h.c.} \tag{2}$$

After the electroweak symmetry breaking, neutrinos get tiny Majorana masses:

$$m_\nu = m_{\nu_D}^T M_R^{-1} m_{\nu_D}, \tag{3}$$

where $m_{\nu_D} = f_\nu \langle H_u \rangle$. The neutrino mixing matrix V_{MNS} is defined by

$$V_{MNS}^T m_\nu V_{MNS} = \text{diag.}(m_{\nu_1}, \ m_{\nu_2}, \ m_{\nu_3}), \tag{4}$$

$$\nu_{F\alpha} = (V_{MNS})_{\alpha i} \nu_{Mi}, \tag{5}$$

where $\nu_{F(M)}$ is a flavor-(mass-) eigenstate of neutrinos. We also define lepton mixing matrices V and V_R, which diagonalize the neutrino Yukawa matrix as follows:

$$V_R^\dagger f_\nu V \equiv f_\nu^{\text{diag}} = \text{diag.}(f_{\nu1}, f_{\nu2}, f_{\nu3}) \tag{6}$$

$$V_{MNS} = VU. \tag{7}$$

Here the matrix U is defined as

$$U^T f_\nu^{\text{diag}} V_R^T M_R^{-1} V_R f_\nu^{\text{diag}} U = \text{diag}(\kappa_1, \kappa_2, \kappa_3). \tag{8}$$

Note that in general the neutrino mixing matrix V_{MNS} is different from the lepton mixing matrix V. As is well known, the neutrino mixing matrix

V_{MNS} is responsible for the neutrino-oscillation physics. The lepton mixing matrix V, on the other hand, is important for the LFV processes, as we will see later. Many models for fermion masses have been proposed so far, in order to accommodate a large mixing for atmospheric neutrinos. It has been pointed out that without introducing a nontrivial structure in the right-handed Majorana neutrino mass matrix, the large mixing for neutrinos and small mixings for quarks are naturally explained by a lopsided structure of the Dirac mass matrices [12, 13]. In this case, the lepton mixing matrix V possesses a large mixing, i.e. $V_{23} \sim O(1)$. As we will see, the experimental limits on the LFV processes can put a stringent constraint on the mixing V.

We should also note that the neutrino mixing V_{MNS} is approximately equal to the mixing matrix V in many models [12, 13, 14]:

$$V_{MNS} \simeq V. \tag{9}$$

In this class of models, neutrino oscillations and LFV phenomena are correlated. Therefore, neutrino experiments, together with the LFV searches, will provide a significant constraint on the models.

3 LFV in SUSY models with right-handed neutrinos

Let us now discuss the LFV in the SUSY models with right-handed neutrinos [7, 8]. Non-zero neutrino Yukawa couplings f_ν generate the LFV in left-handed slepton masses via the renormalization group (RG) effects, even if a common SUSY-breaking mass for all scalars is assumed at the unification (cut-off) scale M_*. An approximate solution to the RG equation for masses responsible for LFV at the weak scale is given by

$$
\begin{aligned}
(\Delta m_{\tilde{L}}^2)_{ij} &\simeq -\frac{(6 + a_0^2)m_0^2}{16\pi^2}(f_\nu^\dagger f_\nu)_{ij} \log \frac{M_*}{M_R}, \\
&= -\frac{(6 + a_0^2)m_0^2}{16\pi^2} V_{ik} V_{jk}^* |f_{\nu k}|^2 \log \frac{M_*}{M_R}, \quad \text{for } i \neq j.
\end{aligned}
\tag{10}
$$

Here, we have assumed a common SUSY-breaking mass (m_0) for all scalar bosons and a common A-term ($A_f = a_0 m_0 f_f$) at the unification scale ($M_* = 3 \times 10^{16}$ GeV), and we have used Eq. (6). Note that the LFV masses $(\Delta m_{\tilde{L}}^2)_{ij}$ depend on the lepton mixing matrix V rather than on the neutrino mixing

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V_{MNS}. LFV processes $\mu \to e\gamma$ and $\mu \to e$ conversion in nuclei are induced by $(\Delta m_{\tilde{L}}^2)_{21}$ component through the slepton-mediated diagrams. Assuming a hierarchical structure of the neutrino Yukawa couplings $|f_{\nu 3}| \gg |f_{\nu 2}| \gg |f_{\nu 1}|$, which is similar to those for the charged-leptons and quarks, the dominant contribution is given by

$$(\Delta m_{\tilde{L}}^2)_{21} \simeq -\frac{(6 + a_0^2) m_0^2}{16\pi^2} V_{23} V_{13}^* |f_{\nu 3}|^2 \log \frac{M_*}{M_R} + \cdots. \qquad (11)$$

A branching ratio for $\mu \to e\gamma$ decay is

$$\mathrm{Br}(\mu \to e\gamma) \simeq F \left| \frac{(\Delta m_{\tilde{L}}^2)_{21}}{(500 \text{ GeV})^2} \right|^2 \tan^2 \beta, \qquad (12)$$

where F is a complicated function of the SUSY parameters. See Refs. [8, 17] for details.

In the models with the Yukawa-coupling unification, the τ neutrino Yukawa coupling is unified into the top-quark Yukawa coupling at M_*. Since the large neutrino Yukawa coupling induces large LFV masses in the slepton sector [8], the event rates for LFV processes can be significantly large. As one can see in Eq. (11), the event rates for $\mu \to e\gamma$ process depend on V_{23} and V_{13} components of the lepton mixing matrix V. The atmospheric neutrino results indicate that the lepton sector has a large mixing between the second and third generations. Especially if the large mixing for atmospheric neutrinos originates from the lepton mixing matrix V, the component V_{23} has a nearly maximal mixing, $V_{23} \sim 1/\sqrt{2}$, and hence it further enhances the LFV. For example, the neutrino-mass models with lopsided Froggatt–Nielsen (FN) U(1) charges in Ref. [13] possess a large mixing in V_{23}, as shown in Fig.1. In many of the existing models [12, 13, 14], the almost maximal mixing comes from the lepton mixing V_{23} as listed in Table 1. Therefore, the searches for LFV can either constrain, or unveil, such a large lepton mixing, $V_{13}V_{23}$, which cannot be probed by neutrino-oscillation physics.

In our analysis, we numerically solve the RG equations and use the complete formula in Ref. [8] for a calculation of the $\mu \to e\gamma$ branching ratio. We fix the heaviest neutrino mass to be 5×10^{-2} eV in order to determine the right-handed neutrino mass scale ($M_R \sim 2 \times 10^{14}$ GeV).

In Fig. 2, we present our numerical result, which shows an upper bound on $V_{13}V_{23}$ from the current limit on the branching ratio for $\mu \to e\gamma$, $\mathrm{Br}(\mu \to e\gamma) < 1.2 \times 10^{-11}$, assuming a Wino mass ($M_2$) to be 150 GeV, $\tan \beta = 3$,
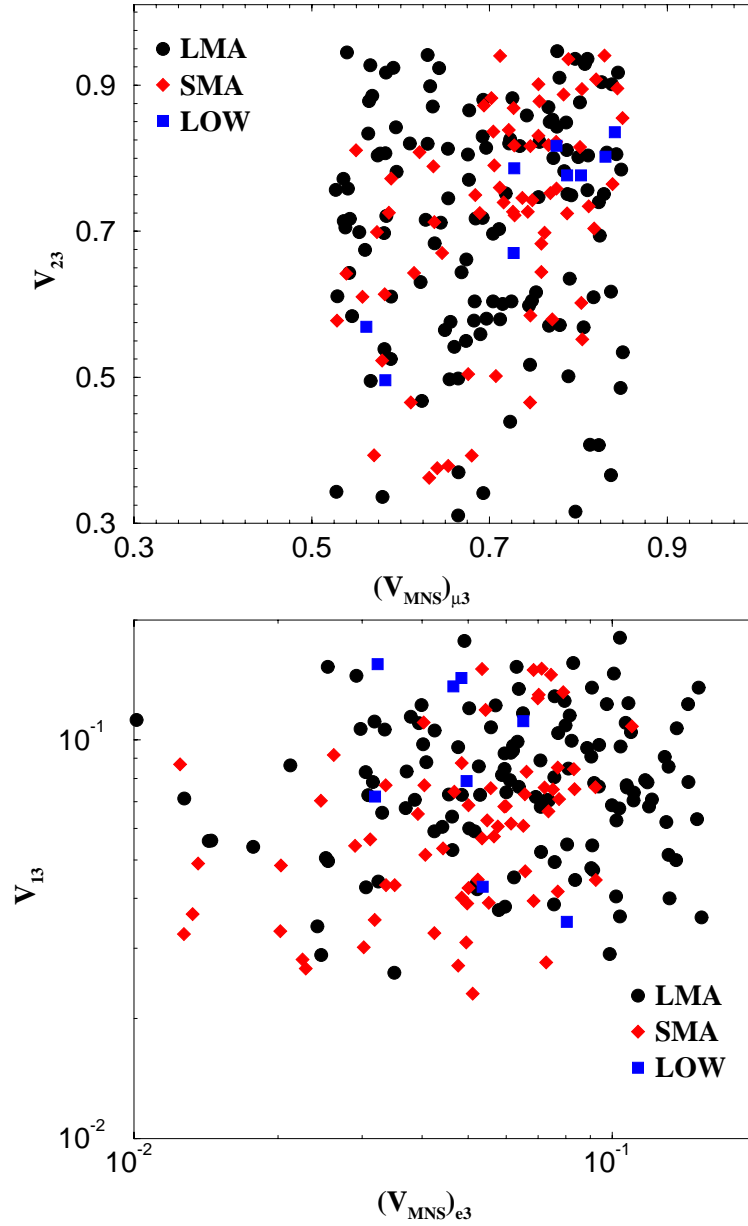
Figure 1: Predicted values of $(V_{MNS})_{\mu 3}$ *vs.* V_{23} and $(V_{MNS})_{e3}$ *vs.* V_{13} in the neutrino-mass models with lopsided Frogatt–Nielsen (FN) U(1) charges in Ref. [13]. Here we considered the FN models where the left-handed lepton doublets L_i ($i = 1$–3) have lopsided FN charges, 1, 0, 0, respectively, while the right-handed charged leptons $\bar{E}_i$ ($i = 1$–3) and the right-handed neutrinos $\bar{N}_i$ ($i = 1$–3) have the same FN charges, 2, 1, 0, respectively. Points denoted by LMA, SMA, and LOW are MSW large mixing angle, MSW small mixing angle, and LOW solutions for solar neutrinos, respectively. For more details, see our forthcoming paper [19].

Models in Refs. [12, 13, 14]	V_{23}	V_{13}	$(V_{MNS})_{e3}$
Albright *et al.* [14]	0.9	0.06	0.05
Altarelli *et al.* [14]	0.5	0.09	0.06
Bando *et al.* [14]	~ 0.7	~ 0.1	~ 0.1
Hagiwara *et al.* [14]	0.7	0.06	0.06
Nomura *et al.* [14]	0.7	~ 0.1	~ 0.1
Sato *et al.* and Buchmüller *et al.* [12, 13]	0.7	~ 0.05	~ 0.05

Table 1: Typical predicted values for V_{23}, V_{13}, and $(V_{MNS})_{e3}$ in various models [12, 13, 14, 18].

and $f_{\nu 3} = f_{\rm top}$, where $f_{\rm top}$ is the Yukawa coupling for the top quark at the unification scale M_*. Here, we do not impose the exact bottom-tau Yukawa unification at M_*, but a milder unification ($f_\tau \simeq f_b$ with 20% deviations) [20] is adopted.[1]

In order to obtain a conservative bound on $V_{13}V_{23}$, we have neglected sub-dominant contributions in the LFV slepton masses in Eq. (11), which are proportional to $|f_{\nu 2}|^2$ and $|f_{\nu 1}|^2$.[2] Moreover, since the branching ratio for $\mu \to e\gamma$ is approximately proportional to $\tan^2 \beta$, we took a small $\tan \beta$ ($\tan \beta = 3$) in Fig. 2. Therefore, our result in Fig. 2 should be considered as a very conservative one.

Our result is also applicable to GUT models with the Yukawa-coupling unification $f_{\nu 3} = f_{\rm top}$. In the GUT models, the RG running from the Planck scale to the unification scale also induces the LFV in slepton masses. Therefore the constraint would be more stringent in general, unless there are accidental cancellations between the two contributions below and above the GUT scale, although the GUT contribution depends on a detail of the model [22].

As can be seen in Fig. 2, the current limit on the $\mathrm{Br}(\mu \to e\gamma)$ can put a severe bound on $V_{13}V_{23}$; typically the limit is $V_{13} < 0.02$ for $V_{23} = 1/\sqrt{2}$ and $m_{\tilde{e}_L} < 500$ GeV. Without any symmetry, such a small value of V_{13} would be very unnatural. Actually as can been seen in Fig. 1 and Table 1, many of the existing models are already strongly constrained or excluded.

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[1]If the exact bottom-tau Yukawa coupling unification is imposed, we need a large $\tan \beta$ ($\tan \beta \sim 50$) [21]. In this case, the constraint on $V_{13}V_{23}$ is more stringent, since the event rate for $\mu \to e\gamma$ is nearly proportional to $\tan^2 \beta$.

[2]If V_{13} is equal to zero, even the sub-dominant contributions are important. For details, see Ref. [17].

Furthermore, the future $\mu \to e\gamma$ experiment [9] with a sensitivity of 10^{-14} in the branching ratio will bring down the limit of V_{13} to 8×10^{-4} for $V_{23} = 1/\sqrt{2}$ and $m_{\tilde{e}_L} < 500$ GeV, and hence it will be able to test the models with Yukawa-coupling unification. In addition to the $\mu \to e\gamma$ process, $\mu \to e$ conversion in nuclei is also important. A ratio between the branching ratio for $\mu \to e\gamma$ and the $\mu \to e$ conversion rate is given by

$$\frac{\mathrm{R}(\mu \to e \text{ in Ti (Al))}}{\mathrm{Br}(\mu \to e\gamma)} \simeq 5 \ (3) \times 10^{-3}, \tag{13}$$

in almost the entire parameter space of the models [8]. Therefore, the future MECO experiment [10] for $\mu \to e$ conversion in Al, with a sensitivity of 10^{-16}, and a further future project, with a sensitivity of 10^{-18} (*e.g.* PRISM [11] for $\mu \to e$ conversion in Ti) will also provide a robust probe on the models; otherwise the LFV phenomena will be observed.[3]

We should stress that even if V_{23} is smaller than the maximal mixing $1/\sqrt{2}$, the constraint does not change much unless it is extremely small. For example, a factor of 2 smaller value of V_{23} gives only a factor of 2 weaker limit on V_{13}.

Finally, we comment on a connection between neutrino oscillation and LFV. As we mentioned in the previous section, it is very likely in many models that the neutrino mixing matrix V_{MNS} equals the lepton mixing matrix V approximately (see Eq. (9)). Therefore, a precise measurement of neutrino mixing $(V_{MNS})_{e3}$ in future neutrino experiments [15, 16] would be very important, since the observation of non-zero $(V_{MNS})_{e3}$ could be in conflict with the constraint from the LFV processes. Therefore, the neutrino oscillation experiments, together with the LFV searches, have a strong potential to exclude a large class of SUSY standard models with the Yukawa-coupling unification.

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[3]A search for another LFV process $\tau \to \mu\gamma$, where the event rate depends on the different matrix elements (V_{23} and V_{33}), gives a different information on the parameters (see Ref. [19]).
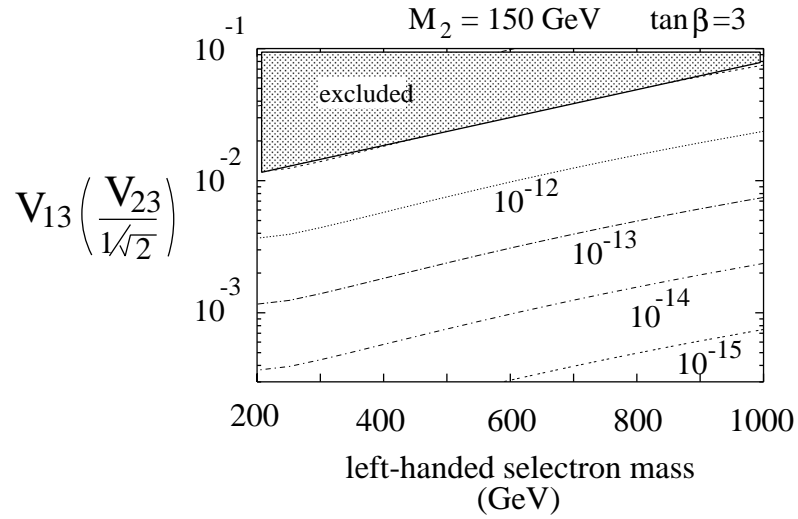
Figure 2: Upper limit on $V_{13}V_{23}$ from the present bound on the branching ratio for $\mu \rightarrow e\gamma$ process (Br($\mu \rightarrow e\gamma$)$< 1.2 \times 10^{-11}$) as a function of the left-handed selectron mass. Here V_{23} is normalized by $1/\sqrt{2}$, and we assume the Wino mass (M_2) to be 150 GeV, and $\tan\beta = 3$. Numbers on the figure denote Br($\mu \rightarrow e\gamma$).

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