No scale Sugra inflation with Type-I seesaw

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Abstract

We show that MSSM with three right handed neutrinos incorporating a renormalizable Type-I seesaw superpotential and no-scale SURGA Kähler potential can lead to a Starobinsky kind of inflation potential along a flat direction associated with gauge invariant combination of Higgs, slepton and right handed sneutrino superfields. The inflation conditions put constraints on the Dirac Yukawa coupling and the Majorana masses required for the neutrino masses and also demands the tuning among the parameters. The scale of inflation is set by the mass of the heaviest right handed neutrino. We also fit the neutrino data from oscillation experiments at low scale using the effective RGEs of MSSM with three right handed neutrinos.

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I. INTRODUCTION

The standard model (SM) extended with three right handed neutrinos (RHNs) is an appealing mechanism for explaining the light neutrino masses through seesaw mechanism \[1-6\]. The SM extended with supersymmetry (SUSY) \[7\] can stabilize the electroweak vacuum and help in achieving the gauge coupling unification. Extensions of SM is also motivated by the need to explain matter anti-matter asymmetry and dark matter. In addition, the theory of inflation, the most successful mechanism for explaining the large scale structure and anisotropy in the CMB spectrum \[8-11\] also needs extension of SM. The CMB observations \[8-11\], in particular, the bounds on the tensor to scalar ratio, \(r_{0.05} < 0.07\) at 95% CL and spectral index, \(n_s = 0.968 \pm 0.006\), put stringent constraints on generic models of inflation like those arising from the quartic potential and quadratic scalar potentials. The \(R + \alpha R^2\) inflation model of Starobinsky\[12\], on the other hand, successfully survives these stringent constraints on \(r\) and \(n_s\) with the prediction that \(n_s = 1 - \frac{2}{N}\) and \(r = 12/N^2 \sim 0.002 - 0.004\) with the minimum number of e-foldings, \(N \sim 55\).

In \[13\], it has been shown that the Starobinsky potential for inflation can be derived from supergravity (SURGA) with a no-scale \[14-16\] Kähler potential and a Wess Zumino superpotential with specific couplings. The realization of the Starobinsky inflation in various GUT models like SO(10) \[17-20\], SU(5) \[21, 22\] and flipped SU(5)xU(1) \[23\] have been studied.

In this work, we consider the minimal supersymmetric standard model (MSSM) extended with three right handed neutrinos with no-scale Kähler potential. The gauge invariant combination of a left handed sneutrino, the SM Higgs and a right handed sneutrino along a D-flat direction acts as inflaton. We show that this leads to Starobinsky type inflationary potential \[12\]. For realization of the Starobinsky potential in this scenario, the Yukawa coupling is related to the Majorana mass (of the heaviest right handed neutrino) which determines the inflation dynamics. This parameter is fixed by the observed amplitude of the temperature anisotropies. This value of right handed neutrino mass is then used as an input for the see-saw neutrino mass generation. We use the RGEs of MSSM with three RHN \[24\] to run down the Yukawa and eliminate the three right handed neutrinos at their respective scale to calculate the effective dimension five operator \(\kappa\) \[25\]. Then, we run the effective dimension five parameter \(\kappa\) using it’s RGEs \[26, 29\] to \(M_Z\) (mass of Z boson) to fit
the neutrino data (mass square differences and three mixing angles).

The inflation with right handed sneutrino has been studied in [30–37] and sneutrino-Higgs along flat direction in MSSM has been studied in [38–53]. In particularly, [52] also considers the gauge invariant combination of a left handed sneutrino, the SM Higgs and a right handed sneutrino along a D-flat direction, but the approach is different. The SUSY breaking scale is $\sim 10^{13}$ GeV which also sets the scale of inflation. The neutrino is a Dirac fermion and extremely tiny $O(10^{-12})$ third generation Yukawa coupling is considered. In our case neutrino is a Majorana particle and the neutrino Yukawa are of same order as SM Yukawa. The scale of inflation is set by the mass of third generation right handed neutrino ($10^{13}$ GeV) and SUSY breaking scale is $O(50-100 \text{ TeV})$.

In present work, the SUSY breaking is done by adding an additional Polonyi field which can acquire vacuum expectation value (vev) at the end of inflation. The mass of the Polonyi field is more than the mass of gravitino by choosing particular superpotential parameters to evade Polonyi problem and to get small cosmological constant [54–58]. After the end of inflation, the decay of RHN can explain the asymmetry through leptogenesis which can be converted into the Baryon asymmetry through Sphaleron process [59].

In Section 2, we present our inflation model. In Section 3, we briefly discuss about reheating and SUSY breaking in our model. In Section 4, we give a benchmark input satisfying the inflationary conditions along with the neutrino oscillation data. Then, we conclude with a brief discussion.

II. INFLATION ALONG D-FLAT LNH DIRECTION

The superpotential considered for inflation contains the terms sufficient to give neutrino masses via type-I seesaw mechanism. The relevant superpotential is given as,

$$W = Y_{\nu}^{ij} L_i H_u N_j + \frac{1}{2} M_N^{ij} N_j N_j + \mu H_u H_d$$

First term is the Dirac term and the second term is Majorana mass term for right handed neutrinos ($N$). Here, $i, j$ represents the number of generations of fermions. Also, $Y_{\nu}$ is a complex matrix and $M_N$ is real diagonal. Here, $H_d$ is another Higgs doublet required for anomaly cancellation in MSSM. The Kähler potential is assumed to be of the general form
given as,

$$K = -3 \ln \left( T + T^* - \sum_i k_i \phi_i \phi_i^* \right)$$  \hspace{1cm} (2)

where we will choose the constants $k_i = 1/3$ for the fields $L_1, N_3, H_u$ whose linear combination will constitute the inflaton and we will choose $k_i \ll 1$ for all other fields (we are working in the units where Planck mass scale, $M_P = (8\pi G)^{-1} = 1$). Therefore for the inflation components the Kähler potential is of the no-scale SURGA form, given as,

$$K = -3 \ln \left( T + T^* - \frac{1}{3}(|L_1|^2 + |N_3|^2 + |H_u|^2) \right)$$  \hspace{1cm} (3)

and for all other fields the Kähler potential goes to the canonical form

$$K = \delta_{ij} \phi_i \phi_j^*$$  \hspace{1cm} (4)

The corresponding potential is given as,

$$V = e^G \left[ \frac{\partial G}{\partial \phi^m} K^m_{\phi} \frac{\partial G}{\partial \phi^*_n} - 3 \right] + \frac{1}{2} D^a D^a$$  \hspace{1cm} (5)

where,

$$G = K + \ln W + \ln W^* \hspace{1cm} D^a = g^a (K_m (T^a)_m \phi_n)$$  \hspace{1cm} (6)

and $K^m_{\phi}$ is the inverse of Kähler metric $K^\phi_n$. Here m, n runs over the number of fields, $g$ is the gauge coupling and $T^a$'s are the generators of each gauge group in SM. The kinetic term is given as $K^m_{\phi} \partial \phi^m \partial \phi^*_n$.

The scalar part of $SU(2)_L$ Higgs and lepton doublets appearing in the Yukawa term during inflation (setting charged component to zero) can be written as,

$$L_i = \begin{pmatrix} \tilde{\nu}_i \\ 0 \end{pmatrix} ; \hspace{1cm} H_u = \begin{pmatrix} 0 \\ h_u \end{pmatrix} ; \hspace{1cm} H_d = \begin{pmatrix} h_d \\ 0 \end{pmatrix}.$$  \hspace{1cm} (7)

The D-flat direction for the gauge invariant combination $LHN$ ($D^a = 0$) is given by,

$$\sum_i |\tilde{\nu}_i|^2 = |h_u|^2$$  \hspace{1cm} (8)

The right handed neutrino is gauge singlet and doesn’t contribute to D-term. Also, we have three generations of neutrinos and the freedom to choose any generation of sneutrino

\footnote{The moduli field $T$ can be stabilised by adding extra terms $((T + T^*)^4 + d(T - T^*)^4)/\Lambda^2$ inside the log term of the Kähler potential \cite{60}. The inflationary potential has a flat direction and Starobinsky form along Re(T)=1 and Im(T)=0 in Planck units (see Fig. (2) of \cite{60}).}
for our inflaton. We chose the 3rd generation right handed neutrino assuming the normal hierarchy of neutrino masses and first generation left handed neutrino. We will assume that all other scalar vevs are zero during the course of inflation to make sure that the fields are non zero in $LNH$ flat direction only and are zero in any other flat direction. So, we consider the fields $\tilde{N}_3$, $\tilde{\nu}_1$ and the Higgs field $H_u$ parametrized in terms of a D-flat direction associated with the gauge invariant $LHN$ and $NN$ terms in the superpotential,

$$\tilde{N}_3 = \tilde{\nu}_1 = h_u = \phi,$$

(9)
to be the inflaton. The superpotential and Kähler potential relevant for inflation is given as,

$$W = Y_{\nu}^{13} \phi^3 + M_{N}^{33} \phi^2 \; ; \; \; K = -3 \ln(T + T^* - |\phi|^2)$$

(10)

After simplifying, the potential and the K.E. has the following form,

$$V = \frac{1}{(1 - |\phi|^2)^2} \left| \frac{\partial W}{\partial \phi} \right|^2 \; ; \; \; L_{K.E.} = \frac{3}{(1 - |\phi|^2)^2} \left| \partial^\mu \phi \right|^2$$

(11)

Here, we have assumed that the non-perturbative Planck scale dynamics fixes the value of $T = T^* = \frac{1}{2}$. After fixing the vev for $T$ the kinetic terms of $T$ can be neglected. To get the canonical K.E. terms, we need to redefine our fields in terms of the new field $\chi$ as,

$$\phi = \tanh \frac{\chi}{\sqrt{3}}$$

(12)

Now, for $\chi = x + iy$, the complex part of $\chi$ is fixed to zero during inflation since it has a mass greater than the Hubble rate $[13]$. Then considering the condition $Y_{\nu}^{13} = -M_{N}^{33}$, the potential for the real part of $\chi$ field looks like,

$$V = M_{N}^{332} (1 - e^{-\frac{2x}{\sqrt{3}}})^2$$

(13)

This is the Starobinsky kind of potential for inflation. Here the scale of the inflation is set by the mass of the heaviest right handed neutrino mass, $M_{N}^{33}$. All other fields have conventional $m^2 \phi^2$ potentials which subdominant ($m^2 \ll M_{N}^{33}$) and steep compared to the inflation potential and so that the fields stole at zero during the slow roll of the inflation without destabilizing the inflation potential. The slow roll parameters for this potential are given by,

$$\eta = -\frac{8e^{-\frac{2x}{\sqrt{3}}}}{3 \left(1 - e^{-\frac{2x}{\sqrt{3}}} \right)^2}; \; \; \; \epsilon = \frac{8e^{-\frac{2y}{\sqrt{3}}}}{3 \left(1 - e^{-\frac{2y}{\sqrt{3}}} \right)^2}.$$

(14)
When $\eta \approx 1$, inflation ends and this corresponds to field value, $x^{\text{end}} \approx .5$. The required number of $N_{e-\text{folds}}=55$ to have sufficient inflation gives the initial field value of $x \approx 4.35$. The power spectrum for scalar perturbation,

$$P_R = \frac{V}{24\pi^2\epsilon} = \frac{M_{33}^{32} \sinh^4 \left( \frac{x}{\sqrt{3}} \right)}{4\pi^2},$$

(15)

requires the value of $M_{33}^{33} = 7.87 \times 10^{-6}$ in Planck units for the central value of $P_R = 2.2 \times 10^{-9}$ given by Planck data [10]. The spectral index $n_s = 0.964$ and tensor to scalar perturbation ratio, $r = 0.002$ for $N_{e-\text{folds}}=55$. However, the deviation from the condition $Y_{\nu}^{13} = -M_{33}^{33}$, even at forth decimal place leads to large deviation from the observational data [13]. The soft mass terms of the inflaton fields will also contribute to the inflation potential. But, we take SUSY breaking scale O(TeV) small as compared to $M_{33}^{33}$. We discuss about SUSY breaking in the next section.

III. REHEATING AND SUSY BREAKING

After the end of inflation, the hot big bang conditions can be restored when energy stored in the inflaton is converted into a thermal bath of the MSSM degrees of freedom. The time required to thermalise the inflaton energy and the resulting reheat temperature $T_{rh}$ depends on the post inflationary dynamics of the LHN flat direction [61]. In [61], the post inflationary dynamics of the LHN flat direction has been studied in $U(1)_R \times U(1)_{B-L}$. Here, we haven’t extended our gauge sector, but the procedure follows the same. Due to the Yukawa couplings and gauge coupling of inflaton fields (however, $N$ has only Yukawa coupling), the inflaton energy is likely to decay rapidly (within one Hubble time) through the so-called instant preheating mechanism [62, 63] to radiation bath of MSSM degrees of freedom. So in present scenario, the estimate for the maximum reheat temperature is given as [61],

$$T_{rh} = \left( \frac{30}{\pi^2 g_*} \right)^{\frac{1}{4}} V_0^{\frac{1}{4}} \sim 10^{13} GeV$$

(16)

Here, $g_* = 228.75$ is the MSSM degrees of freedom. In this model, one can explain the baryon asymmetry, $n_B/n_\gamma$, through leptogenesis. The large reheat temperature ensures the thermal production of the heavy neutrinos $N_3$ whose decay could produce leptogenesis. However
any existing lepton asymmetry will be washed out by the Higgs-neutrino scattering upto
the temperature \( T_{\text{washout}} = 10^{12} \text{ GeV} \). The lepton asymmetry arising from the decay of
the lighter right handed neutrinos \( N_1 \) and \( N_2 \) whose masses are less than \( 10^{12} \text{ GeV} \) can in
principle generate leptogenesis which can be converted to baryogenesis by spahlerons.

Also, an important point to be made is that such a large reheating temperature can pro-
duce relativistic populations of gravitinos, which are dangerous if the lifetimes of gravitinos
are larger than the time of nucleosynthesis, \( \tau_N \sim 1 \text{ sec} \), and their decay after nucleosynthe-
sis will overcome the entire universe. This is famous “gravitino problem” and the general
solution of this problem is to have the graviton mass sufficiently large so that it decay before
nucleosynthesis [64].

\[ \tau_{\text{grav}} \sim 10^5 \text{sec} \left( \frac{1 \text{TeV}}{m_{3/2}} \right)^3 \ll \tau_N \sim 1 \text{sec} \] (17)

So, the viability of this model requires supersymmetry breaking scale \( O(\text{TeV}) \) with gravitino
mass \( m_{3/2} \sim 50 \text{ TeV} \). However, with gravitino mass \( O(50) \text{ TeV} \), there is an upper bound
on the reheating temperature \( \sim O(10^9) \text{ GeV} \) [65]. For \( T_{rh} > 10^9 \text{ GeV} \), the lightest super-
symmetric particle (LSP) produced from the decay of gravitino may over-close the universe.
The possible way out is that the overproduced LSP should decay through small R-partity
violation before BBN [66 67].

The minimal superpotential and Kähler potential responsible for inflation can not give
rise to SUSY breaking. The SUSY can be broken by adding a Polonyi field, \( S \) [57 58] and
adding the following terms,

\[ K(S, \bar{S}) = S\bar{S} + \frac{(S\bar{S})^2}{\Lambda^2} \; ; \; \quad W(S) = M^2 S + \Delta \] (18)

to the Kähler potential and Superpotential respectively.

The term, \( (S\bar{S})^2/\Lambda^2 \) with \( \Lambda \ll 1 \) and the fine tuning of the constant, \( \Delta \) help in the strong
stabilization of the Polonyi field and fixing the vanishingly small cosmological constant \( \sim
10^{-120} \). Also, the cosmological Polonyi problem [54 56] can be solved with the condition,

\[ m_S^2 \gg m_{3/2}^2, \] (19)

so that, it decay into gravitinos. This can be achieved with \( \Delta \neq 0 \) and for \( \Lambda \ll 1 \) and the
potential minimum \( V_{\text{min}} \approx -3\Delta^2 + M^4 \) with \( S_{\text{min}} \approx \Delta \Lambda^2/2M^2 \). Therefore, for \( M^2 \approx \sqrt{3}\Delta, \)
the cosmological constant is very small $\sim 10^{-120}$. This gives $S_{\text{min}} \approx \Lambda^2/2\sqrt{3}$. The gravitino mass is given as

$$m_{3/2}^2 = e^G = \frac{1}{(T + T^*)^3} \ln (SS + \frac{(SS)^2}{\Lambda^2}) |W(S)|^2.$$  \hspace{1cm} (20)

So, at the minimum of the potential the garvitino and Polonyi field masses (in Planck units) are obtained as,

$$m_{3/2}^2 = \Delta^2, \quad m_S^2 = \frac{12\Delta^2}{\Lambda^2} = \frac{12m_{3/2}^2}{\Lambda^2} \gg m_{3/2}^2;$$  \hspace{1cm} (21)

respectively. For $\Lambda \sim 10^{-2}$ and $\Delta \simeq 2 \times 10^{-15}$, we obtain $m_{3/2} \sim 50$ TeV and $m_S \sim 500$ TeV. The inflation potential after SUSY breaking takes the form

$$V = M_{33}^3 \left(1 - e^{-\frac{2x}{\sqrt{3}}} \right)^2 + m_{3/2}^2 \left(\tanh \frac{x}{\sqrt{3}} \right)^2.$$  \hspace{1cm} (22)

With $m_{3/2} = 50$ TeV $\ll M_{33}$ the second term is subdominant compared to the first term in (22). In addition, the slow roll parameters are thus same as that given in 14. The SUSY breaking terms, therefore, does not destabilize the inflaton potential.

IV. EXAMPLE FIT TO NEUTRINO OSCILLATION DATA

The superpotential given in Eqn. (1) is also responsible for neutrino masses through Type-I seesaw. The required ingredients for Type-I seesaw are $Y_{ij}^\nu$ and $M_{ij}^N$. The mass matrix, $M$ for the $\nu$ and $N$ from eqn. (1) can be written as,

$$M = \begin{pmatrix} 0 & M_D \\ M_D^T & M_N \end{pmatrix}.$$  \hspace{1cm} (23)

Here, $M_D = Y_{\nu}v$ with $v=246$ GeV, the SM $v_{ee}$. For $M_N \gg M_D$, the masses of right handed neutrinos are given by $M_N$ and the tiny masses of left handed neutrinos are given as,

$$M_\nu = \frac{1}{2} M_D^T M_N^{-1} M_D.$$  \hspace{1cm} (24)

However, the three right handed neutrinos masses are not degenerate, so we use the RGEs of Yukawa of MSSM with three RHN [24] and eliminates the three right handed neutrinos at their thresholds to calculate the effective dimension five operator, $\kappa$ [25]. Then, we run this effective dimension five parameter, $\kappa$ using RGEs [26,29] to $M_Z$ and calculate the neutrino mass matrix ($3 \times 3$) at $M_Z$. It is given as,

$$M_\nu = v^2 \kappa(M_Z)$$  \hspace{1cm} (25)
where $\kappa = \frac{1}{2} Y_\nu M_N^{-1} Y_\nu$. Using this $M_\nu$, we fit the Neutrino oscillation data along with satisfying the inflation conditions. For this, we use the standard parametrization of the PMNS matrix given by,

$$
U_\nu = \begin{pmatrix}
c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\
-c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\
s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13}
\end{pmatrix} P \quad (26)
$$

where $c_{ij} = \cos \theta_{ij}$, $s_{ij} = \sin \theta_{ij}$ and the phase matrix $P = \text{diag}(1, e^{i\phi_1}, e^{i(\phi_2+\delta)})$ contains the Majorana phases.

Also, we have considered the following constraints on neutrino masses from various experiments.

- The Planck 2015 results put an upper limit on the sum of active light neutrino masses to be [10]

$$
\Sigma = m_1 + m_2 + m_3 < 0.23 \text{ eV}. \quad (27)
$$

The global analysis [68, 69] of neutrino oscillation measurements with three light active neutrinos give the oscillation parameters in their $3\sigma$ range, for normal hierarchy (NH) for which $m_3 > m_2 > m_1$:

- Mass squared differences

$$
\Delta m^2_{21}/10^{-5}\text{eV}^2 = (7.03 \rightarrow 8.09)
$$

$$
\Delta m^2_{31}/10^{-3}\text{eV}^2 = (2.407 \rightarrow 2.643) \quad (28)
$$

- Mixing angles

$$
\sin^2 \theta_{12} = (0.271 \rightarrow 0.345)
$$

$$
\sin^2 \theta_{23} = (0.385 \rightarrow 0.635)
$$

$$
\sin^2 \theta_{13} = (0.01934 \rightarrow 0.02392) \quad (29)
$$

- Dirac Phase

$$
\delta_{PMNS} = (0 \rightarrow 2\pi) \quad (30)
$$
We randomly choose the sixteen parameters of complex $3 \times 3$ matrix $Y_\nu$ and two masses of heavy right handed neutrinos while the component $Y_{\nu}^{13}$ is determined in terms of $M_N^{33} = 7.87 \times 10^{-6} \, M_P$ to satisfy the inflation condition given in the inflation section.

We fit to the neutrino oscillation data within $3\sigma$ given by experiment using a downhill simplex method [70]. One example input for $Y_\nu$ and $M_N$ (GeV) is given in eqn. (31) and the corresponding output is given in Table I.

\[
Y_\nu = \begin{pmatrix}
5.21 \times 10^{-6} + 5.59 \times 10^{-6}i & 1.43 \times 10^{-5} - 3.49 \times 10^{-6}i & -7.87 \times 10^{-6} \\
7.72 \times 10^{-4} + 6.98 \times 10^{-6}i & -3.62 \times 10^{-5} + 7.12 \times 10^{-4}i & -4.14 \times 10^{-4} + 2.92 \times 10^{-5}i \\
-2.45 \times 10^{-2} - 8.26 \times 10^{-5}i & -2.98 \times 10^{-2} + 3.49 \times 10^{-4}i & -0.11 + 4.72 \times 10^{-2}i
\end{pmatrix},
\]

\[
M_N = \begin{pmatrix}
1.84 \times 10^5 & 0 & 0 \\
0 & 1.29 \times 10^9 & 0 \\
0 & 0 & 1.91 \times 10^{13}
\end{pmatrix}.
\] (31)

| Parameter | Value             |
|-----------|-------------------|
| $(m_{12}^2)/10^{-5}(eV)^2$ | 7.9261 |
| $(m_{23}^2)/10^{-3}(eV)^2$ | 2.4071 |
| $\sin^2 \theta_{L_{12}}$ | 0.2838 |
| $\sin^2 \theta_{L_{23}}$ | 0.4180 |
| $\sin^2 \theta_{L_{13}}$ | 0.0237 |
| $\delta_{PMNS}$ | 3.0245 |
| $\phi_1, \phi_2$ | 4.7266, 6.2218 |

We can see from the example input the smallness of off-diagonal elements of $Y_\nu$ can be easily achieved while having the diagonal entries $O(0.1)$. So, we can achieve inflation and fit the neutrino oscillation data with very realistic Yukawa coupling in this scenario.

V. DISCUSSIONS

In this paper, we have shown how a renormalizable superpotential responsible for neutrino masses in MSSM +3 RHN with a no-scale type Kähler potential can lead to Starobinsky
type inflationary potential. The scale of the inflation is set by the mass of the heaviest right-handed neutrino and is essentially independent of the supersymmetry-breaking parameters. The inflation constraints the seesaw parameter values (Yukawa and Majorana mass), but the freedom of choosing the generation makes this constraint very light and fitting the neutrino oscillation data can be achieved very easily with realistic Yukawa couplings. The high reheat temperature, $T_{rh} \sim 10^{13}$ GeV requires a gravitino mass $\sim 50$ TeV to remain consistent with nucleosynthesis. This model predicts a tensor-scalar ratio, $r = 0.002$ and any observation in the near future above this value will rule out this model. The predicted masses of the heavy right-handed neutrinos can be tested in future in constructing models of leptogenesis by heavy neutrino decays. Also, SUGRA models with non-canonical Kähler potential, like the no-scale model discussed in this paper, predict relations between observables like the scale of inflation, SUSY breaking scale [71] and the non-gaussianity [72] which may be testable in future observations.

VI. ACKNOWLEDGEMENTS

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