Singling out SO(10) GUT models using recent PTA results

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In this work, we construct promising model-building routes towards SO(10) grand unified theory inflation and examine their ability to explain the recent pulsar timing arrays (PTAs) results hinting at a stochastic gravitational wave (GW) background at nanohertz frequencies. We consider a supersymmetric framework within which the so-called doublet-triplet splitting problem is solved without introducing fine-tuning. Additionally, realistic fermion masses and mixings, gauge coupling unification, and cosmic inflation are incorporated by utilizing superfields with representations no higher than the adjoint representation. Among the three possible scenarios, two of these cases require a single adjoint Higgs field, and do not lead to cosmic strings. In contrast, the third scenario featuring two adjoints, can lead to a network of metastable cosmic strings that generates a GW background contribution compatible with the recent PTA findings and testable by various ongoing and upcoming GW observatories.

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I. INTRODUCTION

Global collaboration among pulsar timing arrays (PTAs) (NANOGrav [1], PPTA [2], EPTA [3], and IPTA [4]) previously revealed evidence of common-spectrum noise at nanohertz frequencies. Recent analysis [5], including CPTA [6], EPTA [7], NANOGrav [8], and PPTA [9], identified spatial correlations (Hellings-Downs effect [10]), providing strong support for a stochastic gravitational wave background (SGWB). Although the mergers of supermassive black hole binaries (SMBHBs) are natural astrophysical sources of the SGWB at nanohertz frequencies, the new data somewhat disfavors SMBHBs in explaining the observed PTA SGWB signal [8,11]. The SGWB spectrum from SMBHBs is well-described by a simple characteristic strain power law with $f^{-2/3}$ [12]. It turns out that this simple power law does not fit well the new PTA data towards the higher-frequency range [8]. Moreover, should the SGWB arise from a population of merging SMBHB systems, it is anticipated to exhibit anisotropy, which depends on the distribution in the local universe as well as the statistical characteristics of the SMBHB population. However, no such evidence of anisotropy has been observed in the 15-yr dataset [13].

Therefore, the SGWB likely points toward new physics beyond the Standard Model (SM). One of the explanations that fits well with the data is a metastable cosmic string network (CSN) [14,15]. In fact, by performing a fit to the recent PTA data and by comparing with respect to the standard interpretation in terms of inspiraling SMBHBs, it has been shown [14] that metastable cosmic strings provide a better fit, resulting in a Bayes factor of order 10. Since such cosmic strings (CSs) can arise from the multistep spontaneous breaking of the symmetry group of a grand unified theory (GUT) after cosmic inflation, this raises the question of what can be learned about GUTs from this finding.

GUTs [16–21], combined with supersymmetry (SUSY), offer an appealing framework for a more fundamental theory beyond the SM of elementary particles. GUTs unify the three fundamental forces of the SM, while SUSY provides a natural solution to the gauge hierarchy problem and a potential weakly interacting dark matter candidate when R-parity or matter-parity ensures its stability. SO(10)-based GUTs are particularly interesting as they unify all SM fermions of each family into a single irreducible 16-dimensional representation. This 16-dimensional representation also includes a SM singlet right-handed neutrino, which, through the type-I seesaw mechanism [22–26], generates tiny masses for the SM neutrinos.

Promising GUT models must satisfy proton decay bounds and achieve successful gauge coupling unification. In SUSY GUT models, the $d = 5$ proton decay operators are induced by color-triplet exchange, necessitating the

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superheavy nature of color-triplet states compared to their doublet partners, known as the doublet-triplet splitting (DTS) problem [27,28]. A desirable GUT model should solve the DTS problem without fine-tuning parameters. Since GUTs generate the Yukawa matrices out of joint GUT operators, leading to constraints on the flavor structure, a further challenge consists in realizing viable fermion masses and mixings.

Cosmic inflation [29–32] that solves the horizon and flatness problems of the standard big bang cosmology, and explains the origin of structure formation of the observable Universe, could have a deep connection to SUSY GUT models. In addition to the similarity of the scales of inflation and gauge coupling unification, inflation is also crucial to dilute away unwanted topological defects [31,33] like monopoles which generically form at some stage of GUT symmetry breaking. Furthermore, supersymmetric theories typically possess many flat directions, providing an attractive framework for realizing inflation. While monopoles have to be diluted by inflation, other topological defects, like (metastable) CSs [34] that form after inflation can leave an observable signature in the SGWB.

In this work, we explore supersymmetric SO(10) GUTs that naturally solve the DTS problem, generate realistic fermion masses, and achieve successful gauge coupling unification and inflation. We focus on lower-dimensional field representations and investigate scenarios with Higgs fields no higher than the adjoint representation. Three promising routes for SO(10) GUT model building are identified; two cases use a single-adjoint Higgs field, while the third scenario requires two copies. In the latter case, the intermediate symmetry contains two Abelian factors crucial for CSN formation. For the first time, we construct a realistic SUSY SO(10) GUT scenario (particularly the third scenario), satisfying the mentioned criteria and leads to metastable CSs capable of explaining the recent PTA results for a stochastic GW background at nanohertz frequencies.

II. SO(10) MODEL BUILDING

Two major guiding principles in building realistic models in our framework are the natural DTS [35,36] (see also [37–46]) and employing smaller-dimensional representations. In achieving this, we utilize 45_H and $16_H + \overline{16}_H$ Higgs representations to break the GUT symmetry down to the SM, which is subsequently broken by 10_H (and possibly by $16_H + \overline{16}_H$). The fundamental representation contains weak-doublet and color-triplet states,

$$\begin{aligned} 10_H &= (2_H + 3_H) + (\bar{2}_H + \bar{3}_H) \\ &= (1,2,1/2) + (3,1,-1/3) + \text{c.c.}. \end{aligned}$$

The vacuum expectation value (VEV) of the adjoint, $\langle 45_H \rangle \propto i\tau_2 \otimes \mathrm{diag}(a_1, a_2, a_3, a_4, a_5)$ that breaks the GUT symmetry is expected to provide superheavy masses to both these components. With this setup, one can construct three classes of models:

(i) a single adjoint Higgs with $\langle 45_H \rangle \propto B - L$ generator,
(ii) a single adjoint Higgs with $\langle 45_H \rangle \propto I_{3R}$ generator,
(iii) two adjoint Higgses, one with $\langle 45_H \rangle \propto B - L$ generator and another with $\langle 45'_H \rangle \propto I_{3R}$ generator.

For each model, the superpotential takes the form,

$$\begin{aligned} W = W_{\text{GUT-breaking}} + \underbrace{W_{\text{Inflation}} + W_{\text{Mixed}}}_{W_{\text{Intermedite-breaking}}} \\ + W_{\text{DTS}} + W_{\text{Yukawa}}, \end{aligned}$$

where terms in $W_{\text{GUT-breaking}}$ and $W_{\text{Intermedite-breaking}}$ lead to a consistent symmetry breaking of the GUT symmetry down to the SM gauge group. Terms in W_{DTS} realize DTS without fine-tuning, and the $W_{\text{Inflation}}$ part of the superpotential leads to an inflationary period.

A. $B - L$ case

The symmetry breaking chain in this scenario is given by

$$\begin{aligned} SO(10) &\xrightarrow[45_H]{M_{\text{GUT}}} SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\ &\xrightarrow[16_H + \overline{16}_H]{M_I} SU(3)_C \times SU(2)_L \times U(1)_Y. \end{aligned}$$

The GUT-scale symmetry breaking is achieved via

$$W_{\text{GUT-breaking}} \supset \frac{m_{45}}{2} \mathrm{Tr}[45_H^2] + \frac{\lambda}{4\Lambda} \mathrm{Tr}[45_H^4], \tag{1}$$

with the VEV $\langle 45_H \rangle \propto i\tau_2 \otimes \mathrm{diag}(a, a, a, 0, 0)$.

Note that breaking the GUT symmetry gives rise to superheavy monopoles that must be inflated away. Therefore inflation must take place after the formation of the monopoles. A straightforward option is to utilize hybrid [47–49] inflation (an alternative option is tribrid inflation [50,51]) at the last intermediate symmetry breaking stage, which we achieve via employing $16_H + \overline{16}_H$ that acquire VEVs[1]in the right-handed neutrino direction. Then the relevant superpotential term contributing to inflation takes the following form:

$$W_{\text{Inflation}} \supset \kappa S(\overline{16}_H 16_H - m_{16}^2), \tag{2}$$

which fixes the magnitude of the VEVs $\langle 16_H \overline{16}_H \rangle = m_{16}^2$. Here, S is a GUT singlet superfield, the scalar component of which plays the role of the inflaton.

[1]As a result, the appearance of automatic R-parity from within the SO(10) group is no longer possible. However, a discrete symmetry, such as a Z_2 symmetry (matter parity), can readily be imposed.

Since 45_H and $16_H + \overline{16}_H$ have component fields that share the same quantum numbers,

$$45_H, 16_H, \overline{16}_H \supset (1,1,1) + (3,2,1/6) + (\bar{3},1,-2/3) + \text{c.c.},$$

to avoid additional would-be Goldstone bosons, which would ruin gauge coupling unification, these fields must have nontrivial mixing terms. The simplest possible interaction term, $\overline{16}_H 45_H 16_H$, is not welcome since it would destabilize the VEV of 45_H from the desired "Dimopoulos-Wilczek form".

To circumvent this issue, we introduce a second copy of spinorial representations, $16'_H + \overline{16}'_H$, which do not acquire a VEV in the right-handed neutrino direction. Then a consistent symmetry breaking without additional would-be Goldstone bosons can be achieved via the addition of the following terms in the superpotential:

$$W_{\text{Mixed}} \supset \overline{16}_H(\lambda_1 45_H + \lambda'_1 1_H)16'_H + \overline{16}'_H(\lambda_2 45_H + \lambda'_2 1'_H)16_H. \tag{3}$$

Here, we introduced the "sliding singlets" $1_H^{(\prime)}$, which are assumed to have no other terms in the superpotential that could fix their VEVs. They are needed to allow for vanishing F-terms corresponding to $16'_H, \overline{16}'_H$.

Concerning DTS, remarkably, the specific VEV structure of the 45_H provides masses to only the color-triplets, while the weak doublets remain massless, schematically

$$10_{1H}\langle 45_H\rangle 10_{2H} = \cancelto{0}{\bar{2}_{1H}2_{2H}} + \cancelto{0}{\bar{2}_{2H}2_{1H}} + \bar{3}_{1H}3_{2H} + \bar{3}_{2H}3_{1H}. \tag{4}$$

However, if only the above term is added to the superpotential, then the low-energy spectrum would contain four light doublets instead of the usual two doublets of the MSSM. This would spoil the successful gauge coupling unification of the MSSM. To avoid extra light states, we allow a direct mass term for 10_{2H}, i.e.,

$$10_{2H}10_{2H} = \bar{2}_{2H}2_{2H} + \bar{3}_{2H}3_{2H}. \tag{5}$$

Then, the terms in the superpotential relevant for providing the masses of the doublets and triplets and naturally realizing their splittings are

$$W_{\text{DTS}} \supset \gamma 10_{1H}45_H 10_{2H} + m_{10}10_{2H}10_{2H}. \tag{6}$$

A crucial remark is in order. Assuming that only 10_{1H} couples to the fermions, the term in Eq. (4) by itself does not induce proton decay. Once the term in Eq. (5) is also introduced, together they allow the proton to decay via color-triplet Higgses, since now an effective mass term linking $\bar{3}_{1H}$ and 3_{1H} can be written down after integrating out $\bar{3}_{2H}$ and 3_{2H}. This can be understood schematically as follows:

$$\bar{3}_{1H}\langle 45_H\rangle 3_{2H} \qquad \bar{3}_{2H}m_{10}3_{2H} \qquad \bar{3}_{2H}\langle 45_H\rangle 3_{1H}\,.$$

With a sufficiently large effective triplet mass $\sim M_{\text{GUT}}^2/m_{10}$, the $d=5$ proton decay is suppressed.

B. I_{3R} case

The symmetry breaking chain in this scenario is given by

$$SO(10) \xrightarrow[45_H]{M_{\text{GUT}}} SU(4)_C \times SU(2)_L \times U(1)_R \xrightarrow[16_H+\overline{16}_H]{M_I} SU(3)_C \times SU(2)_L \times U(1)_Y,$$

which is obtained by $\langle 45_H\rangle \propto i\tau_2 \otimes \text{diag}(0,0,0,b,b)$. Although the $W_{\text{GUT-breaking}}$ and $W_{\text{Intermedite-breaking}}$ parts of the superpotential are identical to the $B-L$ case, W_{DTS} takes a different form, which we discuss in the following.

Due to $\langle 45_H\rangle \propto I_{3R}$, we now have the opposite situation compared to the previous case, namely

$$10_{1H}\langle 45_H\rangle 10_{2H} = \bar{2}_{1H}2_{2H} + \bar{2}_{2H}2_{1H} + \cancelto{0}{\bar{3}_{1H}3_{2H}} + \cancelto{0}{\bar{3}_{2H}3_{1H}}. \tag{7}$$

Therefore, a different strategy must be implemented to obtain light doublets and superheavy color-triplets. By noting that $16'_H \supset \bar{2}'_H$ is a $SU(2)_R$ singlet, and, on the contrary, $16'_H \supset \bar{3}'_H$ resides in a $SU(2)_R$ doublet, one obtains a mass only for the color-triplet, and not for the weak doublet, i.e.,

$$\overline{16}'_H\langle 45_H\rangle 16'_H = \cancelto{0}{\bar{2}'_H 2'_H} + \bar{3}'_H 3'_H\,. \tag{8}$$

If only the above term is included in the superpotential, then a pair of triplets will remain massless in addition to one pair of doublets. To provide large masses to all the color-triplets, we add two more terms

$$W_{\text{DTS}} \supset \lambda_3 \overline{16}'_H 45_H 16'_H + \lambda_4 10_H 16_H 16_H + \lambda_5 10_H \overline{16}_H \overline{16}_H. \tag{9}$$

As for the $d=5$ proton decay, assuming the SM fermion masses are coming from their coupling to the 10_H (i.e., neglecting all contributions from the 16_H), the effective triplet mass m_T is approximately given by

$$m_T = -\frac{\lambda_3\lambda_4\lambda_5\langle 16_H\rangle\langle\overline{16}_H\rangle}{2\lambda_1\lambda_2\langle 45_H\rangle}. \tag{10}$$

Choosing somewhat small λ_1, λ_2 allows having $m_T \gtrsim 10^{19}$ GeV, which is required by proton decay constraints.

C. $B-L$ and I_{3R} case

Depending on the values of the VEVs of the two adjoints, various symmetry breaking chains may arise in this scenario, examples of which are (a) $\langle 45_H \rangle > \langle 45'_H \rangle > \langle 16_H \rangle, \langle \overline{16}_H \rangle$:

$$
\begin{aligned}
SO(10) &\xrightarrow[45_H]{M_{\rm GUT}} SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\
&\xrightarrow[45'_H]{M_I} SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L} \\
&\xrightarrow[16_H+\overline{16}_H]{M_{II}} SU(3)_C \times SU(2)_L \times U(1)_Y,
\end{aligned}
$$

(b) $\langle 45'_H \rangle > \langle 45_H \rangle > \langle 16_H \rangle, \langle \overline{16}_H \rangle$:

$$
\begin{aligned}
SO(10) &\xrightarrow[45'_H]{M_{\rm GUT}} SU(4)_C \times SU(2)_L \times U(1)_R \\
&\xrightarrow[45_H]{M_I} SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L} \\
&\xrightarrow[16_H+\overline{16}_H]{M_{II}} SU(3)_C \times SU(2)_L \times U(1)_Y,
\end{aligned}
$$

(c) $\langle 45_H \rangle = \langle 45'_H \rangle > \langle 16_H \rangle, \langle \overline{16}_H \rangle$:

$$
\begin{aligned}
SO(10) &\xrightarrow[45_H+45'_H]{M_{\rm GUT}} SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L} \\
&\xrightarrow[16_H+\overline{16}_H]{M_I} SU(3)_C \times SU(2)_L \times U(1)_Y.
\end{aligned}
$$

In this scenario, for each of the adjoints, the GUT symmetry breaking superpotential consists of the terms given in Eq. (1). Since $\langle 45_H \rangle$ and $\langle 45'_H \rangle$ break SO(10) to the left-right symmetry and quark-lepton symmetry, respectively, the first and the second break the generators in $(3,2,+1/6)+(3,2,-5/6)+(3,1,2/3)+\text{c.c.}$ and $(3,2,+1/6)+(3,2,-5/6)+(1,1,+1)+\text{c.c.}$, respectively. Consequently, there would be additional massless states. To avoid such massless states, we add the following mixing term in the superpotential,

$$
W_{\text{GUT-breaking}} \supset \frac{\eta}{\Lambda} \mathrm{Tr}[45_H.45_H.45'_H.45'_H]. \tag{11}
$$

As before, one requires nontrivial interactions between the spinorial representations and the adjoints to give masses to the would-be Goldstones. For the two adjoints, we now introduce two sets of additional spinorial representations, $16'_H + \overline{16}'_H$ and $16''_H + \overline{16}''_H$, and add the following terms, such that the VEVs of the adjoints are not destabilized:

$$
\begin{aligned}
W_{\rm Mixed} \supset\, & \overline{16}_H(\lambda_1 45_H + \lambda'_1 1_H)16'_H + \overline{16}'_H(\lambda_2 45_H + \lambda'_2 1'_H)16_H \\
&+ \overline{16}_H(\lambda_3 45'_H + \lambda'_3 1''_H)16''_H \\
&+ \overline{16}''_H(\lambda_4 45'_H + \lambda'_4 1'''_H)16_H.
\end{aligned} \tag{12}
$$

For the DTS, we include the term $10_{1H}45_H10_{2H}$. However, here we can construct an example model which does not lead to proton decay at leading order via $d=5$ operators. To this end, we forbid the direct mass term $10_{2H}10_{2H}$. Instead, we include a higher-dimensional operator, $10_{2H}.45'^2.10_{2H}$, such that an effective triplet mass for 3_{1H} and $\bar{3}_{1H}$ cannot be written down, since,

$$
\begin{aligned}
&10_{2H}45'^2_H 10_{2H} = \\
&\bar{2}_{2H}2_{2H} + \bar{2}_{2H}2_{2H} + \cancelto{0}{\bar{3}_{2H}3_{2H}} + \cancelto{0}{\bar{3}_{2H}3_{2H}}.
\end{aligned} \tag{13}
$$

With the inclusion of the above two terms, still one pair of color-triplets and an additional pair of weak doublets remain massless. We cure this by adding a term of the form $\overline{16}''_H 16'_H$ to the superpotential,

$$
W_{\rm DTS} \supset \gamma_1 10_{1H}45_H10_{2H} + \frac{\gamma_2}{\Lambda}10_{2H}45'^2_H10_{2H} + \omega_{16}\overline{16}''_H16'_H, \tag{14}
$$

that leads to a single pair of light doublets, as desired.

It is important to note that all the scenarios discussed above can successfully reproduce correct charged fermion masses and mixings by incorporating suitable higher-dimensional operators. The light neutrinos acquire masses through the standard type-I seesaw mechanism. The Majorana masses for the right-handed neutrinos are generated by the following higher-dimensional operator:

$$
W_{\rm Yukawa} \supset Y_R 16_i 16_j \frac{\overline{16}_H \overline{16}_H}{\Lambda} \sim Y_R \frac{v_R^2}{\Lambda} \nu^c \nu^c. \tag{15}
$$

III. GRAVITATIONAL WAVE SIGNALS

In some of the models we consider, breaking e.g., a simple group into a subgroup that contains a $U(1)$ factor leads to monopole creation. To prevent overclosing the universe, inflation must get rid of the monopoles. At some later stage, once the left-over Abelian symmetry is broken, strings appear (we assume the ideal Nambu-Goto string approximation, where the dominant radiation emission of CSs is into GWs [52]). If these two scales are very close, Schwinger nucleation of monopole-antimonopole pairs [53–55] on the string cuts it into pieces and makes it decay. How quickly these metastable strings decay depends on a parameter κ_m [56],

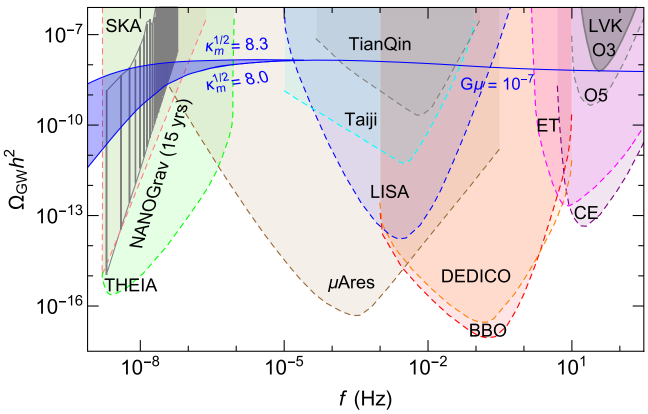

FIG. 1. Examples of GW signals (which we estimate following [57]) from metastable cosmic strings explaining the recent PTA result [14] while being at the edge of the sensitivity reach of LVK [58,59]. The plot also shows examples of sensitivities of possible future observatories (SKA [60], THEIA [61], μAres [62], LISA [63], Taiji [64], TianQin [65], BBO [66], DECIGO [67], ET [68], and CE [69]) which can test the signal at various frequencies. Moreover, the recent NANOGrav 15 yrs result is also shown with gray lines. f is the GW frequency observed today.

$$\kappa_m = \frac{m^2}{\mu} \sim \frac{8\pi}{g^2}\left(\frac{v_m}{v_R}\right)^2, \tag{16}$$

where m is the mass of the monopole and v_m (v_R) is the monopole (string) creation scale. The network behaves like a stable-string network for $\kappa_m^{1/2} \gg 10$.

Metastable CSNs provide an intriguing explanation for the newly released PTA data [14]. The data indicates string tension (μ) values in the range $G\mu \sim 10^{-8}$–10^{-5} for $\kappa_m^{1/2} \sim 7.7$–8.3 (with a strong correlation, cf. Fig. 10 of [14]), consistent with CMB bounds. Notably, the 68% credible region in the $G\mu - \kappa_m^{1/2}$ parameter plane overlaps with the third advanced LIGO-Virgo-KAGRA (LVK) bound, while major parts of the 95% credible region are compatible, preferring $G\mu \lesssim 10^{-7}$ and $\kappa_m^{1/2} \sim 8$ [14], as shown in Fig. 1. From this figure, it is evident that SGWB originating from cosmic strings has a characteristic spectrum that spans over many orders of magnitude in frequencies and can be probed in several GW observatories. This specific feature clearly distinguishes a scenario like ours from other possible explanations for the recent PTA data, such as first-order phase transitions, annihilating domain walls, etc..

However, it should be remarked that the computation of the GW spectrum from metastable CSs carries significant uncertainty [70]. Furthermore, various possible effects are not included in the above shown GW spectrum, for instance, an extended matter domination phase after inflation [70–72] or the change of degrees of freedom below the SUSY breaking scale [71]. Nevertheless, observing a higher-frequency SGWB signal in the next LIGO-Virgo-KAGRA rounds would be a fascinating confirmation of the scenario.

Interestingly, $G\mu \sim 10^{-7}$ corresponds roughly to $v_R \sim 10^{15}$ GeV, which is fully consistent with the type-I seesaw contribution to neutrino masses and corresponds to the right scale for inflation. On the other hand, stable CSs are disfavored by the recent PTA data.[2]

The first (and second) model studied, the $B-L$ (and I_{3R}) case, leads to embedded strings, which are generally unstable [82–84]. Interestingly, all three models in the $B-L$ and I_{3R} case have the potential to produce metastable strings for nearly degenerate monopole and string formation scales: $M_I \sim M_{II}$ for cases (a) and (b), and $M_{\rm GUT} \sim M_I$ for case (c). However, in case (c), a lower GUT scale $\sim 10^{15}$ GeV would have to be arranged that requires suppression of $d=6$ proton decay utilizing the freedom in the Yukawa sector, which makes this case somewhat less appealing. We like to point out that the class of promising SO(10) models we considered in this work may or may not lead to the formation of CSs, contrary to the class of models considered in [85], where the appearance of CSs is unavoidable.

Before concluding, we discuss the gauge coupling unification for an example scenario that leads to metastable CSs [specifically, we choose case (a) within $B-L$ and I_{3R}]. To achieve metastable strings, the monopole and string formation scales must nearly coincide. Therefore, we effectively have three scales; the GUT scale, the monopole/string formation scale, and the SUSY breaking scale (fixed at 3 TeV). To simplify the analysis, we assume that the fields breaking a symmetry are degenerate with the corresponding scale, while the remaining states have GUT-scale masses. This minimal number of free parameters allows us to find a wide range for the monopole/string formation scale, approximately $M_I \sim M_{II} \sim [10^9$–$10^{17}]$ GeV (with 10^{16} GeV $\leq M_{\rm GUT} \leq 10^{18}$ GeV and $M_{\rm GUT} > M_I$), while still being consistent with gauge coupling unification. Our analysis considers a 1% uncertainty on the measured values of the gauge couplings to account for GUT threshold uncertainties.

A comprehensive analysis encompassing gauge coupling unification, fermion masses and mixings, proton decay, GW signal, and the mass spectrum of the component fields from the superpotential terms will be presented in a forthcoming publication.

IV. CONCLUSIONS

We explored promising model-building routes for SO(10) GUT inflation in light of the recent PTA results

[2]Stable cosmic strings, however, were consistent with the previous PTA data. For works on GWs, in light of NANOGrav12.5 data, arising from cosmic strings within GUTs, cf., [73–81].

suggesting the presence of a SGWB at nanohertz frequencies. Our investigation focused on a supersymmetric SO(10) framework with small dimensional representations, effectively solving the doublet-triplet splitting problem without fine-tuning. This approach enables realistic fermion masses, gauge coupling unification, and simple options for embedding cosmic inflation. Among the three model classes studied, one involves two adjoint fields capable of generating a network of metastable cosmic strings. This network generates a SGWB background contribution that can explain the recent PTA data, and will be tested by various upcoming GW observatories.

Note added. As we were completing this work, several papers appeared that also discussed the impact of the new PTA results on new physics scenarios [86–135].

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