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Evading the pulsar constraints on the cosmic string tension in supergravity inflation

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Abstract

The cosmic string is a useful probe of the early Universe and may give us a clue to physics at high energy scales where any artificial particle accelerators cannot reach. Although one of the most promising tools is the cosmic microwave background, the constraint from gravitational waves is becoming so stringent that one may not hope to detect its signatures in the cosmic microwave background. In this paper, we construct a scenario that contains cosmic strings observable in the cosmic microwave background while evading the constraint imposed by the recent pulsar timing data. We argue that cosmic strings with relatively large tension are allowed by delaying the onset of the scaling regime. We also show that this scenario is naturally realized in the context of chaotic inflation in supergravity, where the phase transition is governed by the Hubble induced mass.

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I. INTRODUCTION

Cosmic strings [1] are line-like topological defects which may have been generated in a phase transition in the early Universe through the Kibble mechanism [2][*1]. Contrary to other topological defects such as monopoles and domain walls, they do not dominate the energy density of the Universe, since they evolve according to the scaling law [4]. At first, they were expected to explain the origin of the primordial density fluctuations that lead to structure formation [5]. Although recent observations of the cosmic microwave background (CMB) have ruled out the possibility that cosmic strings alone are responsible for the primordial density fluctuation, they may still contribute to the CMB temperature fluctuation [6–9] which would be a clue to explore physics at high energy scales.

Thus far, many observations for detecting the cosmic strings through the CMB have been done by using the Gott-Kaiser-Stebbins effect [10]. The geodesics of photons passing near a moving string are perturbed, which affects the temperature fluctuations of the CMB. One can thereby constrain the line energy density or the tension, μ. The latest constraint on the string tension is $G\mu \lesssim 10^{-7}$ [6, 7][*2] with G being the gravitational constant. Future observations such as Planck [11, 12] and CMBpol [13] are expected to detect the signatures of strings on small scales or in larger multipoles of the power spectrum with a tension up to $G\mu \sim 10^{-8}$ [8, 9].

There is yet another probe to detect or constrain cosmic strings, the stochastic gravitational wave background (GWB) generated by oscillating string loops [14–19]. When infinite cosmic strings in a Hubble volume intersect with each other, they reconnect and generate closed string loops. As a consequence, although the number of infinite strings in a Hubble volume would increase as the Hubble parameter decreases, the distribution of infinite strings actually obeys the scaling rule [20–22]. At the same time, the loops produced in such a way oscillate, emit gravitational waves (GWs), and shrink gradually. Since the frequency of GWs is determined by the size of cosmic string loops that are generated continuously in the cosmic history, the spectrum of GWB, which are the sum of GWs emitted by each loop, are expected to range widely, from 10^{-10} Hz to 10^{5} Hz or even higher frequencies [14].

[*1] See [3], however, for the impossibility of thermal phase transitions to produce strings relevant to structure formation even in the classical big bang cosmology.

[*2] To be precise, the constraint on the string tension depends on the model of string evolution.

Although their amplitude is typically very small, we may observe their signatures in the GWB through the future gravitational detectors [19] such as eLISA [23], DECIGO [24], and BBO [25] as well as the ongoing pulsar timing arrays (PTAs) [26, 27]. In PTAs, GWs can be detected through the modulation of arrival of pulses from the millisecond pulsars (MSPs) due to the change of the distance between the Earth and MSPs caused by the GWB.

While the PTAs have not detected the signatures of strings thus far, they have obtained stringent constraints on the cosmic string tension [26, 27]. In particular, European Pulsar Timing Array (EPTA) has reported the most stringent constraint, $G\mu < 10^{-9}$, recently [27], though it depends on some assumptions on the string loop formation and GW emission[*3]. Therefore, if we believe this constraint, we cannot hope to detect the signature of cosmic strings in the CMB.

In this paper, we reconsider the formation mechanism of cosmic strings and propose a new scenario which does not contradict with the present GW observations while accommodating large enough $G\mu$ to be detectable in the CMB. If they enter the scaling regime in a later epoch, the loop formation are reduced and hence the amplitude of GWs are suppressed. This can be realized if the phase transition takes place during inflation as proposed in several literatures [3, 29–34]. While previous proposals assumed nontrivial interactions between the string-forming Higgs field and the inflaton field [30–32] or gravity [3, 33], we here point out that the F-term inflation in supergravity can naturally realizes phase transition during inflation thanks to the Hubble induced mass [35] as studied in Refs. [36, 37]. To be concrete, we construct a model that generates cosmic stings during chaotic inflation [38] in supergravity [39]. In this model, the separation of infinite strings is expanded so large that they cannot make loops until long after inflation and hence they enter the scaling regime at a later epoch, for example, around the matter-radiation equality. Since the sources of GWs do not exist at early times, we can evade the PTA constraint on the GWB.

The rest of this paper is organized as follows. In section II, we briefly review the GWs emitted by cosmic string loops and see the current limit from observation using MSPs. In section III, we propose a new scenario of generation and evolution of cosmic strings which

[*3] Recently, Ref. [28] claims a mild constraint, $G\mu < 5 \times 10^{-7}$, based on the observation of EPTA [27]. Although it gives the most conservative constraint, it applies only for a particular value of the model parameter quantifying the size of the loops, which is disfavored by recent numerical simulations [22]. Therefore, we do not adopt it here.

have consistency with this current observation and possibility of detection in the future. In section IV, we construct a model realizing such a scenario, and consider how to determine the parameters of our model. Section V is devoted to discussion.

II. GRAVITATIONAL WAVES FROM STRING LOOPS AND PTA LIMIT

Here we review the gravitational wave emission from cosmic string loops. In order to estimate the GWB today, first we have to know the size distribution of loops at a given time. This can be obtained by considering the scaling rule of strings. Without any intersection and reconnection, the energy density of long string network would fall in proportion to a^{-2}. Actually, however, these processes frequently occur in cosmic time scale to transfer a part of its energy to loops which decay radiating gravitational waves. As a result, the system relaxes to the scaling solution and the energy density of long strings decreases in proportion to $d_{\rm H}^{-2} = H^2$. In terms of the characteristic length of the strings, $\xi = \gamma d_{\rm H} = \gamma H^{-1}$, where γ is a numerical coefficient, we can express the energy density of long strings as

$$\rho = \frac{\xi\mu}{\xi^3} = \frac{\mu}{\xi^2}\,. \tag{1}$$

The coefficient γ is obtained by numerical simulation. Recent simulation suggests $\gamma = 0.15$ in the radiation dominant era, and $\gamma = 0.17$ in the matter dominant era [22] .

Now let us determine the loop distribution function $n(t, \ell)$, which represents the number density of loops with size ℓ at time t. At the moment of loop formation ($t = t_g$), its size is proportional to the horizon size. We can take the initial length as αt_g with α a constant. The length of a loop becomes smaller with time due to the loss of energy by gravitational wave emission. The energy transferred from a loop into gravitational waves per unit time is expressed as $\Gamma G\mu^2$, where Γ is a constant around 50 [21]. Thus, the length of a loop at time t generated at t_g can be written as

$$\ell = \alpha t_g - \Gamma G\mu(t - t_g)\,. \tag{2}$$

In other words, loops with size ℓ at t were formed at

$$t_g = \frac{\ell + \Gamma G\mu t}{\alpha + \Gamma G\mu}\,, \tag{3}$$

while those with size $\ell + d\ell$ at t were formed at

$$t_g + dt_g = \frac{\ell + d\ell + \Gamma G\mu t}{\alpha + \Gamma G\mu}\,. \tag{4}$$

In the flat FRW universe dominated by a perfect fluid with its equation of state given by $P = w\rho$, the scale factor evolves as $a \propto t^{\frac{2}{3(1+w)}}$, and the Hubble parameter is given by $H = \dfrac{2}{3(1+w)t}$. Taking these facts into account, we obtain the following functional form.

$$
\begin{aligned}
n(t,\ell) &= \frac{8(1+3w)}{27(1+w)^3\gamma^2 t_g^3} \times \frac{1}{\alpha(\ell+\Gamma G\mu t)} \times \left(\frac{a(t_g)}{a(t)}\right)^3 \\
&\approx \frac{10(\alpha+\Gamma G\mu)^3}{\alpha(\ell+\Gamma G\mu t)^4} \times \left(\frac{a(t_g)}{a(t)}\right)^3 . \qquad (5)
\end{aligned}
$$

The last approximation is valid for loops which were generated in both the radiation- ($w = 1/3$) and matter-dominated ($w = 0$) regimes .

Now we briefly summarize the formulation to estimate the energy density of GWB. This has been done by Damour and Vilenkin [15], and some modifications have been done in Ref. [18], whose notation we follow. Taking the effect of cusps into account, the Fourier mode of the amplitude, h, of gravitational radiation emitted by a loop with length ℓ at redshift z is given by

$$
h(f,z,\ell) = \frac{G\mu H_0\, \ell^{\frac{2}{3}}}{(1+z)^{\frac{1}{3}}\varphi_r(z)} |f|^{-\frac{4}{3}} \Theta[1-\theta_m(f,z,\ell)] \,. \qquad (6)
$$

Here $\varphi_r(z)$ represents the comoving distance to the loop defined as

$$
r = \frac{1}{H_0}\int_0^z \frac{dz'}{\sqrt{\Omega_\Lambda + \Omega_{m,0}(1+z)^3 + \Omega_{r,0}(1+z)^4}} \equiv \frac{1}{H_0}\varphi_r(z) \,, \qquad (7)
$$

and $\theta_{\rm m}$ is defined by $\theta_{\rm m} \equiv \left(\frac{(1+z)f\ell}{2}\right)^{-\frac{1}{3}}$. The argument of the Heaviside function Θ corresponds to the lowest frequency of loop oscillation. The energy density of GWB is obtained by integrating the contribution from various loops,

$$
\begin{aligned}
\Omega_{\rm GW}(f) &= \frac{4\pi^2}{3H_0^2} f^3 \int dz \int dl\, h^2(f,z,\ell) \frac{d^2R(z,\ell)}{dz d\ell} \\
&= \frac{4\pi^2}{3H_0^2} f^3 \int dz \int dh\, h^2 \frac{d^2R(z,h)}{dz dh}. \qquad (8)
\end{aligned}
$$

Here $\dfrac{d^2R(z,\ell)}{dz d\ell}$ is the observable burst rate per length per redshift

$$
\frac{d^2R(z,\ell)}{dz d\ell} = \frac{1}{H_0^3}\varphi_V(z)\frac{1}{1+z}\frac{2n\,(t(z),\ell)}{\ell}\frac{\theta_m^2(f,z,\ell)}{4}\Theta[1-\theta_m(f,z,\ell)] \,, \qquad (9)
$$

where $H_0^{-3}\varphi_V(z)dz$ represents the spatial volume corresponding to $z \sim z + dz$. Its detailed derivation is found in Ref. [18].

Next, let us consider the observational constraint on the GWB. Currently the strongest limit comes from the Pulsar Timing Array (PTA) experiment, which aims at the detection of the GWB using MSPs. (Detailed explanation and current status is found in Ref. [27].) This experiment has sensitivity at $f = 10^{-9} \sim 10^{-7}$Hz, which is determined by the duration of the observation. At this very low frequency, considering the effect of cusps, one can conclude that the GWB spectrum of energy density $\Omega_{\rm GW} = \dfrac{1}{\rho_{cr}} \dfrac{d\rho_{\rm GW}}{d\log f}$ is proportional to $f^{-\frac{1}{3}}$ from the loops which are generated and emitted gravitational waves during matter-dominated era.

According to Ref. [27] the latest limit is

$$\Omega_{\rm GW}(f = 1{\rm yr}^{-1})h^2 \leq \begin{cases} 10^{-8.4}(2\sigma), \\ 10^{-9.0}(1\sigma). \end{cases} \tag{10}$$

This can be translated to the upper limit on the line density of cosmic string as [27],

$$G\mu \leq 4.0 \times 10^{-9} \,. \tag{11}$$

Note that this constraint is obtained using the method of Damour and Vikenkin [15] taking $\alpha > \Gamma G\mu$ and incorporating all the relevant modes.[*4] On the other hand, a significantly weaker bound has been obtained in Ref. [28] ,

$$G\mu < 5.3 \times 10^{-7} \,. \tag{12}$$

In this study, the effect of cusps is not taken into account, and the constraint merely shows the most conservative one that is obtained when $\alpha \simeq 10^{-5}$. However, if we take the effect of cusps into account, loops emit not only the lowest mode waves but also higher mode waves, so that the amplitude of GWBs is enhanced. In addition, recent simulations [22] indicate that the value of α is actually much larger than 10^{-5}. That is why we must face the more stringent limit than (12).

We can evade the PTA constraint on $G\mu$ without taking an unnatural value of α if we consider a different mechanism of string formation. Since the dominant contribution to $\Omega_{\rm GW}(f = 1{\rm yr}^{-1})$ comes from the loops which were generated during late radiation and early matter-dominated eras, we can evade the PTA constraint if such loops can be reduced. In this case, we can explain the PTA constraint without setting $G\mu$ so small.

[*4] The formalism developed in Ref. [18] would yield a slightly more stringent constraint on the line density.

III. A NEW SCENARIO OF STRING FORMATION

We propose a new scenario of formation and evolution of cosmic strings which can evade the PTA constraint. Since the initial separation of long cosmic strings are comparable to the horizon size at the formation time, it seems impossible to change the initial separation drastically to reduce the number of loops, in particular, in the case of ordinary cosmic strings formed by the spontaneous symmetry breaking after reheating or at the end of hybrid inflation [40, 41]. The number of loops can be reduced in a scenario where cosmic strings are formed not after inflation but *during* inflation. This type of scenario was also studied by one of us in a different context [3].

Let us see the detail of this scenario. Since the mean separation of strings evolves larger and larger as inflation proceed, they have to wait for a long time to come across. Strings can neither intersect with each other nor form loops until their mean separation falls well below the horizon. In this scenario, we can remove the sources of the GWB (cosmic string loops) at early times, we can satisfy the observational constraints imposed by GWB with a relatively big $G\mu$.

We calculate $\Omega_{\rm GW}$ in such a scenario where loops start to form only after the time $t_{\rm sc}$, when strings begin to satisfy the scaling law. Although some loops may have been formed before $t_{\rm sc}$, such loops would be rare and hence we neglect them here. The calculation can be done as follows. We perform z integral in Eq. (8) to $z_{\rm sc}$, which corresponds to the redshift of $t_{\rm sc}$. The relation between cosmic time t and z is given by

$$t = \frac{1}{H_0}\int_z^\infty \frac{dz'}{(1+z')\sqrt{\Omega_{\Lambda 0}+(1+z')^3\Omega_{\rm m0}+(1+z')^4\Omega_{\rm r0}}} \equiv \frac{1}{H_0}\varphi_t(z), \tag{13}$$

$$z = \varphi_z(H_0 t) \quad \text{(the inverse function of } \varphi_t(z)\text{)}. \tag{14}$$

Some modifications are needed in h integral in Eq. (8) . Shortly after $t_{\rm sc}$, the minimum loop size is $\ell_1 \equiv \alpha t_{\rm sc} - \Gamma G\mu(t_1 - t_{\rm sc})$ at t_1 since very small loops are absent. Besides, Θ-function means that only the strings with $\ell > \frac{2}{f(1+z)} \equiv \ell_2$ contribute to the GWB. Therefore h integral runs not from 0 but $\max(h_1, h_2)$, where h_1 and h_2 correspond to the amplitude of GWs emitted from the strings with the length ℓ_1 and ℓ_2, respectively. The upper limit of h corresponds to the maximum loop size. Although the authors of Ref. [17] insist that one can remove the contribution from rare bursts by introducing the cutoff in h integral, this cutoff has ambiguity. Therefore we do not introduce it and include "all bursts".

Figure 1 shows the density parameter of GWBs from cosmic strings at the frequency $f = 1$ year^{-1} as a function of $z_{\rm sc}$ for various values of α. We can read off the condition $z_{\rm sc} \simeq 1 - 10^3$ should satisfy for each value of α and $G\mu$ to be consistent with the PTA observation. Figure 2 shows the maximum $z_{\rm sc}$ (or the earliest time for strings to start making loops) in order not to contradict with the PTA data for $\alpha = 10^{-1}$. We find an approximate formula of $z_{\rm sc}$ to evade the PTA limit as

$$z_{\rm sc} \lesssim 3 \times 10^3 \left(\frac{G\mu}{10^{-8}}\right)^{-0.94} . \tag{15}$$

Therefore, we still have the possibility to detect the features of cosmic string by the future CMB observations with small enough $z_{\rm sc}$ since the cosmic strings with $G\mu \lesssim 10^{-7}$ are allowed if loop formation is delayed.

IV. COSMIC STRINGS AND INFLATION IN SUPERGRAVITY

We present a new mechanism to realize the above scenario in supergravity. While the previous models introduced a nonminimal coupling with the scalar curvature to the string-forming scalar field [3, 33] or direct coupling with the inflaton [30–32], our model does not require any additional interactions, since it is automatically provided by gravitationally suppressed interaction arising from supergravity.

The F-term potential in supergravity is written as

$$V = e^{\frac{K}{M_{\rm G}{}^2}} \left[(D_i W) {K^{ij}}^{-1} (D_j W)^* - \frac{3}{M_G^2} |W|^2 \right] , \tag{16}$$

where K is the Kähler potential and W is the superpotential. Here we have defined

$$K^{ij} = \frac{\partial^2 K}{\partial \phi_i \partial \phi_j^*} , \tag{17}$$

$$D_i W = \frac{\partial W}{\partial \phi_i} + \frac{1}{M_G^2} \frac{\partial K}{\partial \phi_i} W , \tag{18}$$

and $M_G = 1/\sqrt{8\pi G}$ is the reduced Planck mass. During the standard slow-roll inflation, the energy density ρ is dominated by the potential energy, so the Friedmann equation reads $3H^2 M_G^2 = \rho \cong V$. For $K \ll M_G^2$ we find

$$3M_G^2 H^2 \cong V \cong (D_i W) {K^{ij}}^{-1} (D_j W)^* - \frac{3}{M_G^2} |W|^2 \tag{19}$$

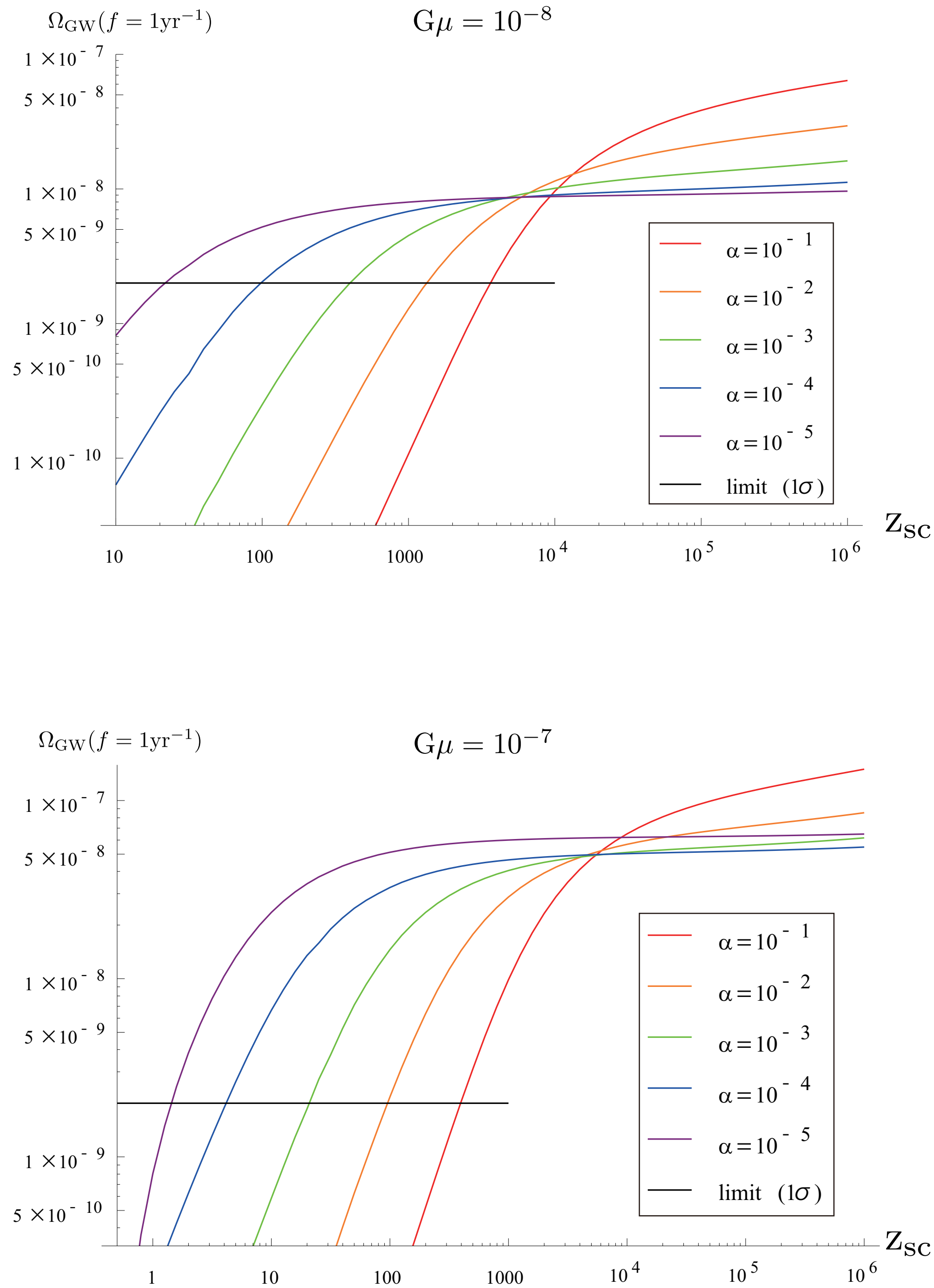

FIG. 1: Energy density of GWB at $f = 1\text{year}^{-1}$, $\Omega_{\rm GW}$, as a function of $z_{\rm sc}$, the redshift at the onset of the scaling solution for various values of α. The horizontal line shows the upper bound imposed by the PTA observation.

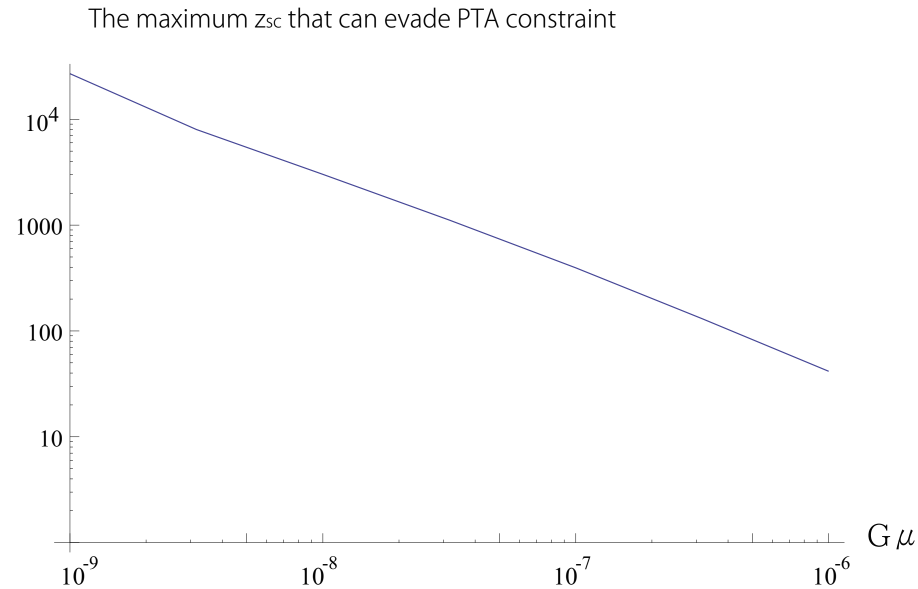

FIG. 2: The maximum $z_{\rm sc}$ that can evade the PTA constraint as a function of $G\mu$ for $\alpha = 10^{-1}$.

in the F-term inflation. Since the kinetic terms of the scalar fields are given by $-K^{ij}\partial_\mu\phi_i\partial^\mu\phi_j^*$, the minimal model assumes $K = |\phi_i|^2$ which yields canonical kinetic terms. Here we can expand the exponential factor after canonical normalization and diagonalization,

$$e^{\frac{K}{{M_{\rm G}}^2}} = 1 + \frac{\phi_i\phi^{i*}}{M_{\rm G}^2} + \cdots . \tag{20}$$

Therefore, the scalar potential V contains the term $3H^2\phi_i\phi^{i*}$. This means that the field ϕ_i receives an additional Hubble-scale mass, which is called the Hubble-induced mass[*5]. If the Hubble parameter is large enough in the beginning of inflation, these fields fall down to the origin quickly and the symmetry is restored during inflation. Moreover, in inflation models in which the Hubble parameter changes gradually during inflation such as the chaotic inflation model [38], it is possible for the phase transition to take place during inflation. Here we

[*5] The inflaton should not acquire the Hubble induced mass; otherwise the slow-roll condition cannot be satisfied. In order to realize inflation, one must assume shift symmetry [39] or adopt contrived models [42].

	ϕ	X	s	ψ	$\bar{\psi}$
$U(1)_S$	0	0	0	$+1$	-1
R	0	$+2$	$+2$	0	0
Z_2	-1	-1	0	0	0

TABLE I: Charge assignments on superfields in the model under the $U(1)_S$ symmetry, the R-symmetry, Z_2-symmetry.

adopt the following superpotential,

$$W = W_I + W_S, \tag{21}$$

$$W_I = M\phi X, \tag{22}$$

$$W_S = m^2 s - \lambda s\psi\bar{\psi}, \tag{23}$$

and the Kähler potential

$$K = \frac{1}{2}(\phi + \phi^*)^2 + |s|^2 + |X|^2 + |\psi|^2 + |\bar{\psi}|^2\,, \tag{24}$$

to realize chaotic inflation, where W_I and W_S are superpotentials of the inflation sector and the string sector, respectively. Here ϕ is the inflaton, X is an additional singlet, ψ and $\bar{\psi}$ are the symmetry-breaking Higgs fields that break $U(1)_S$ symmetry, s is a singlet which destabilizes the Higgs field at the origin, M is the inflaton mass, m is the symmetry breaking scale, and λ is a coupling constant. Note that we impose R-symmetry and discrete Z_2 symmetry in order to suppress all other unwanted interaction such as $X\phi^2, s\phi$, and so on[*6]. Charge assignments on the fields are shown in Table I.

This model can realize chaotic inflation naturally thanks to the shift symmetry, $\phi \to \phi + c$ with c being a real parameter, along the imaginary part of ϕ which plays the role of the

[*6] Note that the term $\epsilon s\phi^2$ in the superpotential cannot be forbidden by these symmetries but it breaks the shift symmetry of ϕ. Since we can expect that the inflaton mass M as an order parameter of the shift symmetry breaking, the coupling constant ϵ would be suppressed enough, say $\epsilon \sim M^2 \sim 10^{-10}$ in Planck units. Thus, following discussion does not change.

inflaton [39]. Setting $\phi = \frac{1}{\sqrt{2}}(\chi + i\varphi)$, the potential for these fields reads

$$
\begin{aligned}
V &= e^{\frac{K}{M_{\rm G}{}^2}} \left[(D_i W) K^{ij^{-1}} (D_j W)^* - \frac{3}{M_G^2} |W|^2 \right] \\
&= \exp\left[\frac{1}{M_{\rm G}^2} \left(\chi^2 + |X|^2 + |s|^2 + |\psi|^2 + |\bar{\psi}|^2 \right) \right] \\
&\quad \times \left[\frac{1}{2} M^2 (\chi^2 + \varphi^2) + M^2 |X|^2 + m^4 - \lambda m^2 (\psi\bar{\psi} + \psi^* \bar{\psi}^*) \right. \\
&\qquad + \lambda^2 |\psi|^2 |\bar{\psi}|^2 + \lambda^2 |s|^2 |\bar{\psi}|^2 + \lambda^2 |s|^2 |\psi|^2 \\
&\qquad \left. + \frac{1}{M_{\rm G}^2} \left\{ \left(\frac{3}{2} M^2 \chi^2 - \frac{1}{2} M^2 \varphi^2 \right) |X|^2 + \cdots \right\} \right] .
\end{aligned}
\tag{25}
$$

Note that we do not have to consider all of these terms because the fields χ, s, X, ψ, and $\bar{\psi}$ have very steep potential due to the factor $e^{\frac{K}{M_{\rm G}{}^2}}$ at the field values larger than M_G. This means they cannot have field values much bigger than $M_{\rm G}$. Since χ and s have larger masses than the Hubble parameter, they are practically fixed at the origin during inflation. Although X has a mass of order of the inflaton mass, its contribution to the total potential energy (or the dynamics of inflation) and the density fluctuation are negligible due to its much smaller field value than the inflaton. Since only the inflaton φ can have a value much larger than M_G, this model reduces to a simple chaotic inflation model, $V = M^2\varphi^2/2$.

Now we investigate the phase transition during inflation. Diagonizing the mass matrix of ψ and $\bar{\psi}$, we obtain the following canonically normalized fields:

$$
\begin{cases}
\psi_1 = \dfrac{1}{\sqrt{2}} (\psi + \overline{\psi}^*) \,, \\
\psi_2 = \dfrac{1}{\sqrt{2}} (\psi - \overline{\psi}^*) \,.
\end{cases}
\tag{26}
$$

The relevant terms that control the dynamics of phase transition during inflation are

$$
\begin{aligned}
V &= \frac{1}{2} M^2 \varphi^2 + \left[\frac{1}{M_G^2} \left(\frac{1}{2} M^2 \varphi^2 + m^4 \right) - \lambda m^2 \right] |\psi_1|^2 + \frac{\lambda^2}{4} |\psi_1|^4 \\
&\approx \frac{1}{2} M^2 \varphi^2 + \left(3H^2 - \lambda m^2 \right) |\psi_1|^2 + \frac{\lambda^2}{4} |\psi_1|^4 \,.
\end{aligned}
\tag{27}
$$

The first term is the potential for the inflaton. The mass term of ψ_1 changes its sign from positive to negative during inflation at $H^2 = \lambda m^2/3$. As a result, ψ_1 may be destabilized and the $U(1)_S$ symmetry is spontaneously broken during inflation, while ψ_2 has always positive mass squared.

Our model contains parameters λ, m, and M. M is related to the primordial power spectrum of curvature fluctuations as

$$\Delta_{\mathcal{R}}^2(k) = \frac{1}{24\pi^2}\frac{V}{M_{\mathrm{G}}^4}\frac{1}{\epsilon}\bigg|_{k=aH} \approx \frac{1}{6\pi^2}\frac{M^2}{M_{\mathrm{G}}^2}\mathcal{N}^2, \tag{28}$$

with $\mathcal{N} \simeq 55$ being the number of e-folds when observable scales exit the horizon. The observed value $\Delta_{\mathcal{R}}^2(k) = 2.4 \times 10^{-9}$[43] determines the inflaton mass to be $M \cong 10^{13}$GeV. We assume λ is of order of unity and m is much smaller than M_{G}.

The line energy density of cosmic strings is related to the vacuum expectation value of the string-generating field as [*7]

$$\mu \cong 2\pi\langle{\psi_1}^2\rangle = 4\pi\frac{m^2}{\lambda}\,. \tag{29}$$

Therefore one can determine the value of m^2/λ by fixing $G\mu$.

The crucial issue is whether strings with $G\mu \geq 10^{-8}$, whose tension is large enough to be detectable by CMB observations [9], can be produced with appropriate density so that they can evade the PTA constraint but are observable using CMB. It is therefore important to study when the phase transition takes place to clarify the mean separation and the correlation length of the string network.

In our scenario, the phase transition is triggered by quantum fluctuation during inflation rather than thermal fluctuations and we can make use of the results of [3, 31, 34, 37] to analyze its properties. Calculations based on the stochastic inlfation method [44, 45] show that the phase of the string-forming field ψ_1 is fixed when the classical potential force surpasses the effect of stochastic quantum fluctuations [34]. This occurs $\Delta\mathcal{N} = \sqrt{c/2}$ e-folds after $\psi_1 = 0$ has become classically unstable with $3H^2 = \lambda m^2$ where $c \equiv \lambda m^2/M^2$ [34, 37]. The number of e-folds from this epoch to the end of inflation is given by

$$\mathcal{N}_D = \frac{1}{2}(c - \sqrt{2c} - 1)\,, \tag{30}$$

which gives the extra dilution absent in the case phase transitions occur after or without inflation. The Hubble parameter at this epoch, which we denote with a suffix f, is given by

$$H_f = \left(\frac{c}{3} - \frac{\sqrt{2c}}{3}\right)^{1/2} M. \tag{31}$$

[*7] To be precise, one should treat the coefficient as a function depending on λ, rather than 2π. However within the range of interesting value of λ in our model, we can approximate this coefficient to 2π[41].

The typical separation between strings at this epoch can be estimated as follows [3, 31]. First choose an arbitrary point $\boldsymbol{x}_0$ with an arbitrary field value of $\psi_1(\boldsymbol{x}_0)$. One can always perform a gauge transformation to make it real there: $\psi_1(\boldsymbol{x}_0) \equiv \psi_{1R}(\boldsymbol{x}_0) + i\psi_{1I}(\boldsymbol{x}_0) = \psi_{1R}(\boldsymbol{x}_0)$. Then since a string is a locus of $\psi_{1R}(\boldsymbol{x}) = \psi_{1I}(\boldsymbol{x}) = 0$, the nearby string segment can be found by searching a point (in fact, a line) with $\psi_{1R}(\boldsymbol{x}) = 0$ on the surface of $\psi_{1I}(\boldsymbol{x}) = 0$. To the lowest-order Taylor expansion, the condition a string exists within a distance r from the point $\boldsymbol{x}_0$ is given by

$$r|\boldsymbol{e}_I \times \nabla\psi_{1R}(\boldsymbol{x}_0)| > |\psi_{1R}(\boldsymbol{x}_0)|, \tag{32}$$

where $\boldsymbol{e}_I \equiv \nabla\psi_{1I}(\boldsymbol{x}_0)/|\nabla\psi_{1I}(\boldsymbol{x}_0)|$ is the unit normal vector to the surface $\psi_{1I} = 0$. Since $\psi_{1R}(\boldsymbol{x}_0)$ and $\nabla\psi_{1R}(\boldsymbol{x}_0)$ are statistically independent, almost Gaussian variables, one can express the probability for the inequality to be satisfied at a distance r, $P(r)$, in terms of the variances $\langle\psi_{1R}^2\rangle \equiv \sigma^2$ and $\langle(\nabla\psi_{1I})^2\rangle/3 = \langle(\partial_x\psi_{1I})^2\rangle \equiv \sigma_{g'}^2$, which will be evaluated later. Without a loss of generality one can take z-axis along $\boldsymbol{e}_I$. Then writing $\nabla\psi_{1R} \equiv (u_x, u_y, u_z)$, $P(r)$ is given by

$$\begin{aligned} P(r) &= \int_{r^2(u_x^2+u_y^2)>\psi_{1R}^2} P(\psi_{1R}, u_x, u_y, u_z) d\psi_{1R} d^3u \\ &= \frac{1}{(2\pi)^2\sigma\sigma_{g'}^3} \int \exp\left(-\frac{\psi_{1R}^2}{2\sigma^2} - \frac{u_x^2+u_y^2+u_z^2}{2\sigma_{g'}^2}\right) d\psi_{1R} d^3u \\ &= \frac{\sigma_{g'} r}{(\sigma^2 + \sigma_{g'}^2 r^2)^{1/2}}, \end{aligned} \tag{33}$$

where $P(\psi_{1R}, u_x, u_y, u_z)$ is the one-point probability distribution function of ψ_{1R} and $\nabla\psi_{1R}$. Thus the typical distance to reach a string, which also gives an estimate of the mean separation of strings, is given by $r \simeq \sigma/\sigma_{g'}$.

Next we calculate the variances at $t = t_f$, which is obtained from the power spectrum of ψ_1 [37]

$$\mathcal{P}_{\psi_1} = |\psi_{1k}(t_f)|^2 \simeq \frac{H^2(t_k)}{2k^3}\left(\frac{S(t_k)}{S(t_f)}\right) \exp\left[-\frac{9}{2M^2}\left(H(t_k) - \frac{2M}{3}\sqrt{c - \frac{1}{2}}\right)^2 + \frac{(2-\sqrt{3})^2 c}{4}\right], \tag{34}$$

where

$$S(t) \equiv \left(\lambda m^2 - \frac{M^2}{2} - \frac{3H^2(t)}{4}\right)^{1/2}, \tag{35}$$

and t_k is the epoch when the comoving mode k exit the horizon. We can evaluate the ratio of the variances as

$$\begin{aligned}\frac{\sigma^2}{\sigma_{g'}^2} &= \frac{\int |\psi_{1k}(t_f)|^2 \frac{d^3k}{(2\pi)^3}}{\int \frac{k^2}{3}|\psi_{1k}(t_f)|^2 \frac{d^3k}{(2\pi)^3}} \equiv \left(\frac{r(c)}{H_f}\right)^2 \\ &\simeq 15\sqrt{2\pi}(x_f - x_d)\exp\left[\frac{15}{2}(x_f - x_d)^2 + \frac{c}{5} - \sqrt{2c} + \frac{2}{5}\right]\frac{1}{H_f^2} \qquad (36)\end{aligned}$$

where

$$x_d \equiv \frac{2}{5}\left(c - \frac{1}{2}\right)^{1/2}, \quad x_f \equiv \left(\frac{c}{3} - \frac{\sqrt{2c}}{3}\right)^{1/2}. \qquad (37)$$

Now we get the typical separation of strings. In addition to the expression (36), we also numerically evaluate $r(c)$ as a function of c and get $r(c) \simeq 400$ at $c = 100$.

Finally the typical separation at the time corresponding to redshift z is given by

$$\begin{aligned}d(z) &= re^{\mathcal{N}_D}\frac{a_{\rm R}}{a_{\rm end}}\frac{a_{\rm eq}}{a_{\rm R}}\frac{a(z)}{a_{\rm eq}} \\ &= 3\times 10^{18}\frac{r(c)}{x_f(c)}\exp\left[\frac{1}{2}(c - \sqrt{2c} - 1)\right]\frac{1}{z+1}\left(\frac{T_{\rm R}}{10^6{\rm GeV}}\right)^{-\frac{1}{3}}{\rm GeV}^{-1}, \qquad (38)\end{aligned}$$

where the suffix end, R, and eq represent the epochs of the end of inflation, reheating and radiation-matter equality, respectively . Note that the scale of the string network at z is given by

$$\xi = \gamma d_{\rm H} = \gamma H^{-1} = \frac{\gamma \times 7\times 10^{41}}{\sqrt{\Omega_{\Lambda 0} + \Omega_{\rm m0}(1+z)^3 + \Omega_{\rm r0}(1+z)^4}}{\rm GeV}^{-1}. \qquad (39)$$

One can then get the relation between model parameres and $z_{\rm sc}$ by equating (38) and (39).

We have shown in Figure 3 the parameter corresponding to $z_{\rm sc} = 1100$ and $d(z=0) = H_0^{-1}$ which means there would be only one string inside our present horizon. We find m should be about 10^{14}GeV, and so use $m_{14} \equiv m/10^{14}$GeV. Therefore c becomes $\frac{\lambda m_{14}^2}{100}$. The horizontal axis is related to the time when strings start to scale and produce loops. According to the results of previous section, there is an upper limit on $z_{\rm sc}$ in order not to contradict with the MSP observation. This upper limit on $z_{\rm sc}$ yields a lower limit on λm_{14}^2. The value of $G\mu$ is proportional to the value of vertical axis.

V. DISCUSSION

In this paper, we have studied the new mechanism to evade the PTA constraints on the cosmic string tension. We have found that the PTA constraints can be evaded if the loop

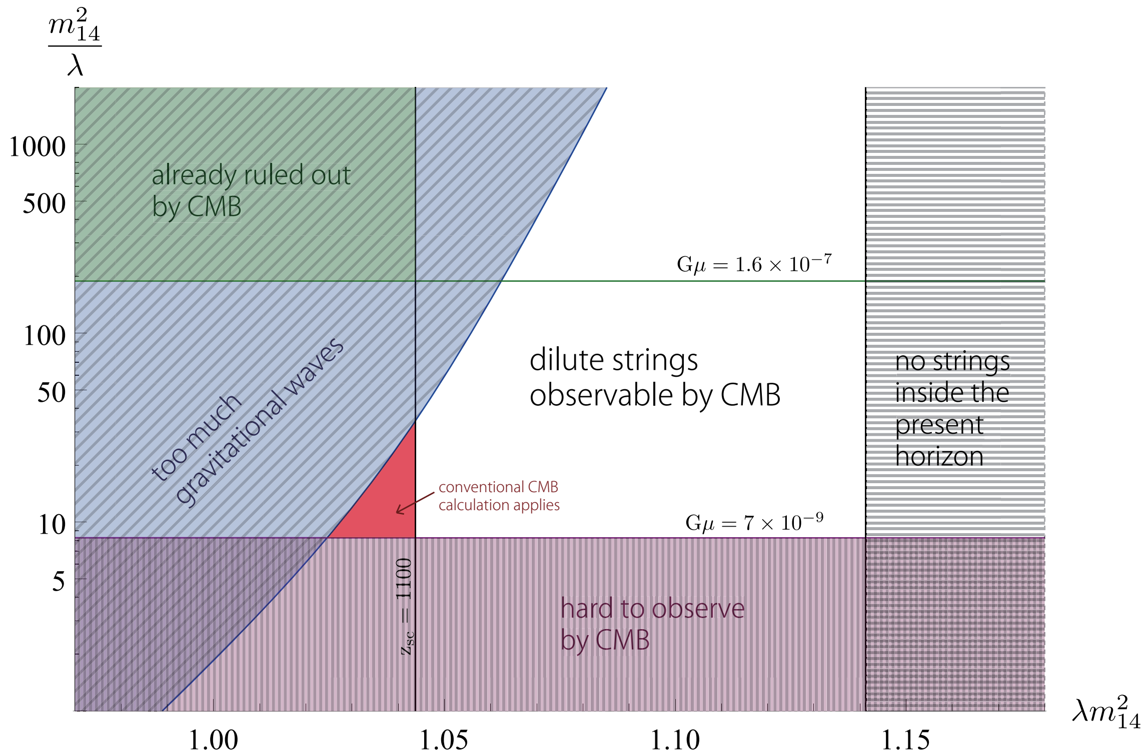

FIG. 3: The allowed parameter region of this model is shown. In Blue-painted region, the network of strings start to produce loops so early that the gravitational waves emitted by them contradict the observation. Though determining this region requires more statistical analysis of bursts and more realistic treatment of α, we emphasize that some region remain allowed. The most interesting region, which represents the case strings can not only evade the PTA constraint but also contribute to the anisotropy of CMB as much as the ordinary strings do. The white region corresponds to the case strings are observable through CMB but its density is smaller than the conventional scenario. As a recent constraint on $G\mu$, we used the value obtained in [7]. According to [9], strings with $G\mu < 7 \times 10^{-9}$ are not observable.

formation starts at a later epoch, say, $z \lesssim 10^4$, depending on the value of α and the tension of the cosmic strings. We have also shown that this scenario can be realized if the phase transition and the string formation take place during inflation. As a specific example, we have shown there is a natural parameter space where we can expect to detect the signatures of cosmic strings through the CMB observation while accommodating the PTA constraints in the F-term chaotic inflation model in supergravity [39].

As seen in Figures 1 the value of α is very important to determine the constraint on

$z_{\rm sc}$. Recent simulation [22] indicates that loops are formed with various sizes ranging from $\alpha \simeq 10^{-1}$ to $\alpha < 10^{-5}$ in the standard scenario, which also shows the inappropriateness of the conclusion of Ref. [28]. In our scenario, however, the mean separation of cosmic strings is much larger than the horizon initially, so it is expected that larger loops are produced more likely than in the ordinary scenario. In addition, inflation also erase small-scale fluctuations on long strings. Hence strings do not develop small-scale structures and the correlation length remains of order of the Hubble radius for some time even after $z_{\rm sc}$. As a result we expect α takes a relatively large value $\alpha \sim 10^{-1}$ at the onset of the scaling regime.

In our scenario, the effect to the CMB may need to be reconsidered. If inflation did not last so long after strings were formed, stings start to satisfy the scaling rule at an earlier epoch. As a consequence, their separation would already become smaller than the horizon at the recombination. Then there would be as many strings as in the standard scenario near the last scattering surface of the CMB photons. The effect to the CMB would be the same as what has already been analyzed in the literature [46]. On the other hand, if inflation lasted very long after strings had been formed, they cannot start to scale until very late time. Therefore, their separation would be larger than the horizon at the recombination. In this case, their effect on the CMB would become smaller. For quantitative understanding, numerical simulations of the string network with appropriate initial conditions are desired.

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