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Gravitational waves as a probe of the SUSY scale

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We investigate the sources of the Hubble-induced mass for a flat direction in supersymmetric theories and show that the sign of the Hubble-induced mass generally changes just after the end of inflation. This implies that global cosmic strings generally form after the end of inflation in a wide class of supersymmetric models, including the minimal supersymmetric Standard Model. The cosmic strings emit gravitational waves whose frequency corresponds to the Hubble scale, until they disappear when the Hubble parameter decreases down to the soft mass of the flat direction. As a result, the peak frequency of gravitational waves is related to the supersymmetric scale. The observation of this gravitational wave signal, in conjunction with results from collider experiments and dark matter searches, can yield information of supersymmetry-breaking parameters.

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I. INTRODUCTION

The observation of gravitational waves (GWs) will open a new window onto the early Universe and provide information on physics at correspondingly high energy scales [1]. Stochastic GW signals are generated by non-equilibrium phenomena during the postinflationary period, such as preheating [2–8], first-order phase transition [9–17], turbulent motions [18–21], topological defects [22–29], and self-ordering scalar fields [30–34]. Quantum fluctuations during inflation are another source of GWs, called the inflationary GW background [35,36]. Each source may predict characteristic GW signals measurable by future GW detectors such as LISA [37], DECIGO [38], Advanced LIGO [39], and ET [40]. The observation of these GW signals will improve our understanding of particle physics beyond the Standard Model.

Supersymmetric (SUSY) theories are well motivated in particle physics, because they address the hierarchy problem and also achieve gauge coupling unification. In SUSY theories, there usually exist scalar fields called flat directions, whose potentials are flat as long as they maintain SUSY and renormalizability (see Ref. [41] for the minimal SUSY Standard Model). In this paper, we investigate the dynamics of a flat direction and show that a cosmic string network generally forms after the end of inflation in a wide class of SUSY models. They eventually disappear when the Hubble parameter decreases down to the mass of the flat direction. Since the cosmic strings generate a stochastic GW background with a peak frequency corresponding to the Hubble scale, the information of the mass of the flat direction is imprinted on the GW spectrum. Combined with collider experiments and dark matter searches, the observation of this GW signal can provide us invaluable hints on the SUSY-breaking scale.

II. DYNAMICS OF FLAT DIRECTION

Let us focus on one flat direction, which we denote as ϕ. During inflation, the flat direction obtains Hubble-induced mass through higher-dimensional Kähler potentials like

$$-\int d^2\theta d^2\bar{\theta}\frac{c'_{H_{\rm inf}}}{M_*^2}|X|^2|\phi|^2, \tag{1}$$

where M_* is a cutoff scale, and $c'_{H_{\rm inf}}$ is an $O(1)$ constant. We expect that M_* is less than or equal to the Planck scale $M_{\rm Pl}$ ($\simeq 2.4\times 10^{18}$ GeV). The F-term of X drives inflation and satisfies the relation of $|F_X|^2 = 3H_{\rm inf}^2 M_{\rm Pl}^2$. This implies that the flat direction obtains a Hubble-induced mass of $c_{H_{\rm inf}}H_{\rm inf}^2|\phi|^2$ during inflation, where

$$c_{H_{\rm inf}} = 3c'_{H_{\rm inf}}\frac{M_{\rm Pl}^2}{M_*^2}. \tag{2}$$

While the coefficient $c_{H_{\rm inf}}$ is assumed to be negative in the context of the Affleck-Dine baryogenesis [42,43], we assume $c_{H_{\rm inf}} > 0$ in this paper. In this case, the flat direction stays at the origin, i.e., $\phi = 0$, during inflation.

After inflation ends and before reheating completes, the energy density of the Universe is dominated by the oscillation of a scalar field (denoted by I) and the Hubble parameter decreases with time as $H(t)\propto a^{-3/2}(t)$, where $a(t)$ is the scale factor. During this oscillation era, the flat direction obtains a Hubble-induced mass through

$$-\int d^2\theta d^2\bar{\theta}\frac{c'_H}{M_*^2}|I|^2|\phi|^2 \supset -\frac{c'_H}{M_*^2}|\dot{I}|^2|\phi|^2. \tag{3}$$

Since the oscillation time scale of I is much smaller than H^{-1}, we can average $|\dot{I}|^2$ over the oscillation time scale like

$|\dot{I}|^2 \approx \rho_I(t)/2 \simeq (3/2)H^2(t)M_{\rm Pl}^2$. Thus, we obtain the Hubble-induced mass of $c_H H^2(t)|\phi|^2$ with

$$c_H = \frac{3}{2} c'_H \frac{M_{\rm Pl}^2}{M_*^2} \tag{4}$$

during the oscillation era. The coefficient of Hubble-induced mass during inflation, $c_{H_{\rm inf}}$, is generally different from the one during the oscillation era, c_H, because the field I is generally different from the field X. For example, in the chaotic inflation model proposed in Ref. [44], the field X (in this paper) is identified with the field X (in Ref. [44]), while I is identified with the inflaton φ. In the simplest hybrid inflation model proposed in Ref. [45] (Ref. [46]), the field X is identified with the field Φ (S), while I is a waterfall field Ψ_1 and Ψ_2 (ϕ and $\bar{\phi}$). While the nonrenormalizable terms are heavily dependent on the high-energy physics beyond the cutoff scale, we assume that among many flat directions in SUSY theories, (at least) one flat direction has $c_{H_{\rm inf}} > 0$ and $c_H < 0$. Let us investigate the dynamics of such a flat direction. The flat direction stays at $\phi = 0$ during inflation, and then it obtains a large VEV after the end of inflation.

If the flat direction has a nonrenormalizable superpotential, the resultant GW background is generally too small to be detected [47]. Thus, we consider the case that the superpotential of the flat direction is absent ($W_\phi = 0$) due to some chiral symmetries, such as discrete R-symmetry [48].[1] In this case, the potential of the flat direction obtains higher-dimensional terms coming from nonrenormalizable Kähler potential and is written as

$$V(\phi) = c_H H^2(t)|\phi|^2 + m_\phi^2|\phi|^2 + a_H H^2(t)\frac{|\phi|^{2n-2}}{M_{\rm Pl}^{2n-4}} + \cdots, \tag{5}$$

where n (≥ 3) is a certain integer, m_ϕ is the soft mass of the flat direction, and a_H is given by

$$a_H = a'_H \left(\frac{M_{\rm Pl}}{M_*}\right)^{2n-2}, \tag{6}$$

with a'_H being an $O(1)$ constant. The dots $\cdots$ represent the other irrelevant higher-dimensional terms. Here we assume that higher-dimensional terms breaking $U(1)$ symmetry like ϕ^n + c.c. are absent due to some symmetries. The potential minimum is determined as

$$\langle\phi\rangle = \left(\frac{|c_H|}{a_H(n-1)}\right)^{1/(2n-4)} M_{\rm Pl} \quad (\sim M_*), \tag{7}$$

as long as $c_H H^2(t) \gg m_\phi^2$.

[1] For another example, $B-L$ symmetry can forbid W_ϕ for, say, $u^c d^c d^c$ flat direction without causing phenomenological problems.

Here we investigate the components of the flat direction in detail. Let us consider the $L_i H_u$ flat direction as an illustration, where i is a flavor index.[2] Since the above nonrenormalizable Kähler potentials generally breaks flavor symmetry, the $L_i H_u$ flat direction has the local $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetry and the global $\mathrm{U}(1)_L$ symmetry. The nonzero VEV of the flat direction breaks the local $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetry to the $\mathrm{U}(1)_{\rm EM}$ symmetry and breaks the global $\mathrm{U}(1)_L$ symmetry completely. As a result, a global cosmic string network forms after the end of inflation.[3] By using numerical simulations (see below), we confirm that the cosmic string network reaches a scaling regime within a certain time. In this regime, the number of cosmic strings in the Hubble volume is $O(1)$.

Since the Hubble-induced mass decreases with time as $\propto H(t) \propto a^{-3/2}(t)$, the soft mass m_ϕ eventually dominates the potential of the flat direction, and the flat direction starts to oscillate around $\phi = 0$. Let us denote that time as $t_{\rm decay}$, which is estimated by

$$c_H H^2(t_{\rm decay}) \simeq m_\phi^2. \tag{8}$$

Note that for low-scale SUSY models we should require $T_{\rm RH} \lesssim 10^9$ GeV to avoid the gravitino problem [49,50]. Even for high-scale SUSY models such as pure gravity mediation, $T_{\rm RH} \lesssim 10^{10}$ GeV is required to avoid an overproduction of LSP [$= O(100)$ GeV] [51–53]. In such well-motivated cases, the cosmic strings disappear before reheating completes—that is, $t_{\rm decay} < t_{\rm RH}$, where $t_{\rm RH}$ is the time at which reheating completes. Hereafter, we consider such a case.

III. CALCULATION OF GWS

GWs are emitted by the dynamics of the cosmic strings [22–25,28,32,33]. We calculate the spectrum of GWs using the method proposed in Refs. [7,27], which is suitable to our situation compared with the method to calculate GW amplitudes from localized sources derived in Ref. [54].[4] Hereafter, we change the time variable from t to the conformal time τ, which is defined by $\mathrm{d}t = a\mathrm{d}\tau$.

[2] The absence of the nonrenormalizable superpotential for $L_i H_u$ is disfavored from the viewpoint of the observed neutrino oscillations. In this paper, however, we take the $L_i H_u$ flat direction as an example to take advantage of its simple flavor structure.

[3] If the flat direction has a larger flavor symmetry like $\mathrm{SU}(3)_{\rm flavor}$, the nonzero VEV of the flat direction breaks the global $\mathrm{SU}(3)_{\rm flavor} \times \mathrm{U}(1)_L$ symmetry to the $\mathrm{SU}(2) \times \mathrm{U}(1)$ symmetry. Also in this case, GWs are emitted by the dynamics of randomly distributed Nambu-Goldstone modes of the flat direction [30–34,47], though cosmic strings are absent.

[4] Since $\mathrm{d}^2V(\phi)/\mathrm{d}\phi^2 \sim H^2(t)$ at $\phi = \langle\phi\rangle$, a typical width of cosmic strings is of the order of the Hubble radius. This implies that the Nambu-Goto approximation, which was used in Refs. [55,56] for example, is inappropriate to describe these cosmic strings.

The energy density of GWs can be calculated from [7]

$$\Omega_{\rm gw}(\tau) \equiv \frac{1}{\rho_{\rm tot}(\tau)} \frac{{\rm d}\rho_{\rm gw}(\tau)}{{\rm d}\log k} \simeq \frac{k^5}{24Va^4H^2} \int {\rm d}\Omega_k \sum_{ij} (|A_{ij}|^2 + |B_{ij}|^2), \quad (9)$$

where $\rho_{\rm tot}(\tau)$ $[= 3M_{\rm Pl}^2 H^2(\tau)]$ is the total energy density of the Universe. The frequency-dependent coefficients A_{ij} and B_{ij} are given as [27]

$$A_{ij}(\mathbf{k}) = -16\pi G \int_{\tau_i}^{\tau_f} {\rm d}\tau' \tau' a(\tau') f_A(k\tau') T_{ij}^{\rm TT}(\tau', \mathbf{k}), \quad (10)$$

$$B_{ij}(\mathbf{k}) = 16\pi G \int_{\tau_i}^{\tau_f} {\rm d}\tau' \tau' a(\tau') f_B(k\tau') T_{ij}^{\rm TT}(\tau', \mathbf{k}), \quad (11)$$

where $T_{ij}^{\rm TT}$ is the Fourier-transformed transverse-traceless part of the anisotropic stress. Here we have implicitly assumed that the source term $T_{ij}^{\rm TT}$ lasts during the interval of $[\tau_i, \tau_f]$. The functions f_A and f_B are calculated by matching solutions of the Einstein equation in the oscillation era with that in the radiation-dominated era at $\tau = \tau_{\rm RH}$ such as [47]

$$f_A(k\tau') = [a_1 n_1(k\tau') - a_2 j_1(k\tau')],$$
$$f_B(k\tau') = [-b_1 n_1(k\tau') + b_2 j_1(k\tau')], \quad (12)$$

where j_l and n_l are the spherical Bessel and Neumann functions of order l, and the coefficients are given as

$$a_1 = x^2[j_1(x)\partial_x n_0(x) - n_0(x)\partial_x j_1(x)]|_{x\to k\tau_{\rm RH}},$$
$$a_2 = x^2[n_1(x)\partial_x n_0(x) - n_0(x)\partial_x n_1(x)]|_{x\to k\tau_{\rm RH}},$$
$$b_1 = -x^2[j_1(x)\partial_x j_0(x) - j_0(x)\partial_x j_1(x)]|_{x\to k\tau_{\rm RH}},$$
$$b_2 = -x^2[n_1(x)\partial_x j_0(x) - j_0(x)\partial_x n_1(x)]|_{x\to k\tau_{\rm RH}}. \quad (13)$$

Here, let us estimate the spectrum of GWs from the cosmic strings before we show our simulation results. GWs are most efficiently emitted for the frequency corresponding to the Hubble scale, which is given as $k \sim \tau^{-1}$. Since the emission of GWs proceeds through the Planck suppressed interaction [see Eqs. (9), (10), and (11)], the produced energy density of GWs can be estimated as[5]

$$\frac{\Delta\Omega_{\rm gw}}{\Delta \log \tau} \sim \left(\frac{\langle\phi\rangle}{M_{\rm Pl}}\right)^4 \sim \left(\frac{M_*}{M_{\rm Pl}}\right)^4, \quad (14)$$

where we use Eq. (7). The GW energy density decreases with time as $\propto a^{-1}(\tau) \propto \tau^{-2}$ in the oscillation era, and the peak frequency of the emitted GWs decreases with time as $\propto \tau^{-1}$. Therefore, the GW spectrum for the frequency $k \gtrsim \tau^{-1}$ is proportional to k^{-2}.[6] For large-scale modes $k \lesssim \tau^{-1}$, the Fourier-transformed transverse-traceless part of the anisotropic stress $T_{ij}^{\rm TT}$ is independent of k due to the loss of causality at the large scale [7,27]. Then, using $j_l(x) \to 2^l l! x^l/(2l+1)!$ and $n_l(x) \to -(2l)!/(2^l l! x^{l+1})$ for $x \ll 1$, we obtain $\Omega_{\rm gw} \propto k$ for $\tau_{\rm RH}^{-1} \ll k \ll \tau^{-1}$ and $\Omega_{\rm gw} \propto k^3$ for $k \ll \tau_{\rm RH}^{-1}$. Thus, the spectrum of GWs bends at the wave number around $k \simeq \tau_{\rm RH}^{-1}$.

The GW emission terminates at the time of $\tau_{\rm decay}$, when the cosmic strings disappear. The GW peak energy density and frequency at that time are roughly estimated as Eq. (14) and $k_{\rm peak} \sim aH(\tau_{\rm decay}) \simeq a(t_{\rm decay}) c_H^{-1/2} m_\phi$, respectively. Then the GW amplitude decreases with time as $\propto a^{-1}(\tau)$ until reheating completes.

One might wonder if (tachyonic) preheating occurs at the beginning and its GW signal affects the one emitted from the cosmic strings. In reality, the GW peak frequency emitted from the preheating is much higher than the one emitted from the cosmic strings at the time they disappear ($\tau = \tau_{\rm decay}$), because the Hubble scale at the time of the preheating is much larger than the one at $\tau = \tau_{\rm decay}$. Since GWs with such high frequencies are beyond the detectability of future experiments, we can neglect GW signals emitted from the preheating. Note also that numerical calculations we have performed include the effect of the preheating.

IV. RESULTS OF NUMERICAL SIMULATIONS

To calculate the GW spectrum from Eqs. (9), (10), and (11), we have performed lattice simulations with $N^3 = 256^3$ grid points in the oscillation era, in which $H(\tau) \propto a^{-3/2}$. We use a numerical method similar to the one used in Ref. [29]. While details are given in Ref. [47], we stress that our results are independent of the choices of the time step, simulation size, and grid size by changing their values by a factor of 50%. Initial fluctuations of the field value are seeded by the vacuum fluctuations, though we have checked that our results are qualitatively insensitive to the detailed form of the initial conditions. Numerical simulations are performed in a unit with $a(\tau_i) \equiv a_i = 1$ and $H(\tau_i) \equiv H_i = 1$, though we explicitly write H_i below.

Figure 1 shows the evolution of GW spectra obtained from our numerical simulations. We confirm that the

[5]If the cutoff scale is equal to the Planck scale ($M_* \simeq M_{\rm Pl}$), the right-hand side of Eq. (14) becomes $O(1)$, which implies that we have to include the effect of the backreaction of GW emission. In addition, the energy density of the cosmic strings is comparable to the energy density of the Universe, and the evolution of the Universe is nontrivial. Hereafter, we consider the case that the cutoff scale is less than the Planck scale.

[6]Although our simulation results (Fig. 1) appear different from k^{-2} for $k \gtrsim \tau^{-1}$, we have confirmed that this is owing to the limitation of the simulation time (see Ref. [47]).

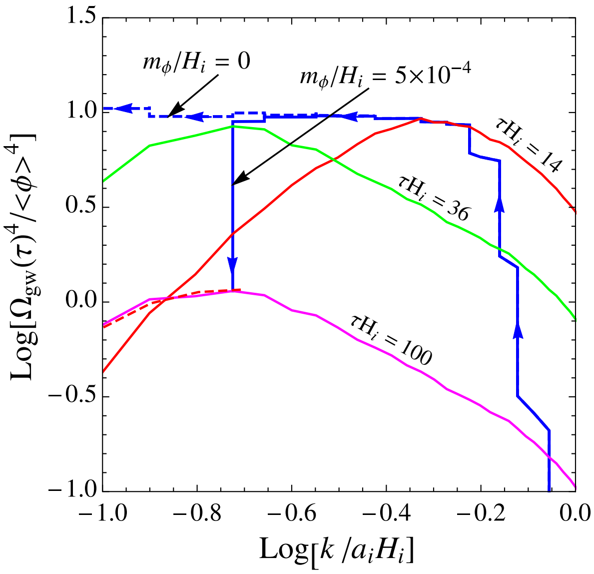

FIG. 1 (color online). Evolution of GW spectra obtained by numerical calculations. We show the obtained spectra at $\tau H_i = 14$ (red line), $\tau H_i = 36$ (green line), and $\tau H_i = 100$ (magenta line). We take $n = 4$, $c_H = 15$, and $m_\phi/H_i = 5\times10^{-4}$. The blue (dashed) curve represents the contour of the peak frequency and the peak GW energy density for the case of $m_\phi/H_i = 5\times10^{-4}$ (0). The red dashed curve represents an analytic estimation given by Eqs. (10) and (11) with k-independent $T_{ij}^{\rm TT}$.

cosmic string network reaches the scaling regime at the time around $\tau H_i \simeq 14$. Then the produced energy density of GWs $\Delta\Omega_{\rm gw}/\Delta\log\tau$ becomes constant while its peak frequency $k_{\rm peak}$ decreases with time as $k_{\rm peak} \propto \tau^{-1}$. When the flat direction starts to oscillate at the time around $\tau H_i = 36$ ($\simeq \tau_{\rm decay} H_i$), the peak frequency becomes constant and the GW spectrum begins to redshift adiabatically as $\propto \tau^{-2}$ ($\propto a^{-1}$). We find that GW spectra for $k \lesssim \tau_{\rm decay}$ are consistent with analytical estimation with $T_{ij}^{\rm TT} = {\rm const.}$ (red dashed curve), which we expect from the viewpoint of causality. From Fig. 1, we obtain numerical factors such that

$$\Omega_{\rm gw}(t_{\rm decay}) \simeq 2\left(\frac{\langle\phi\rangle}{M_{\rm Pl}}\right)^4, \tag{15}$$

$$\frac{k_{\rm peak}}{a(t_{\rm decay})} \simeq 3\frac{m_\phi}{\sqrt{c_H}}, \tag{16}$$

where $\langle\phi\rangle$ and $t_{\rm decay}$ are given by Eqs. (7) and (8), respectively.

V. PRESENT SPECTRUM OF GWS AND DETECTABILITY

Since the GW energy density decreases with time as $\propto a^{-1}(\tau)$ in the oscillation era, i.e., the matter-dominated era, its present value is given as

$$\begin{aligned}\Omega_{\rm gw}h^2(t_0) &\simeq \Omega_r h^2\left(\frac{g_s(t_0)}{g_s(t_{\rm RH})}\right)^{4/3}\left(\frac{g_*(t_{\rm RH})}{g_*(t_0)}\right)\left(\frac{H_{\rm RH}}{H_{\rm decay}}\right)^{2/3}\Omega_{\rm gw}(t_{\rm decay}) \\ &\simeq 2\times10^{-7}\left(\frac{m_\phi}{10^3\ {\rm GeV}}\right)^{-2/3}\left(\frac{T_{\rm RH}}{10^9\ {\rm GeV}}\right)^{4/3}\left(\frac{M_*}{M_{\rm Pl}}\right)^{10/3},\end{aligned} \tag{17}$$

where t_0 is the present time, $\Omega_r h^2$ ($\simeq 4.15\times10^{-5}$) is the present energy density of radiation, and g_* (g_s) is the effective relativistic degrees of freedom for the energy (entropy) density. We have used Eq. (15) and assumed $c'_H = a'_H = 1$ and $n = 4$ in the last line [see Eqs. (4), (6), and (7)]. Taking redshift into account, we obtain the present value of peak frequency f_0 like

$$\begin{aligned}f_0 &\simeq \left(\frac{g_s(t_0)}{g_s(t_{\rm RH})}\right)^{1/3}\left(\frac{T_0}{T_{\rm RH}}\right)\left(\frac{H_{\rm RH}}{H_{\rm decay}}\right)^{2/3}\frac{k_{\rm peak}}{2\pi a(t_{\rm decay})} \\ &\simeq 7\times10^2\ {\rm Hz}\left(\frac{m_\phi}{10^3\ {\rm GeV}}\right)^{1/3}\left(\frac{T_{\rm RH}}{10^9\ {\rm GeV}}\right)^{1/3}\left(\frac{M_*}{M_{\rm Pl}}\right)^{1/3},\end{aligned} \tag{18}$$

where we use Eq. (16) in the last line. As explained above, the spectrum of GWs bends at the wave number corresponding to the Hubble scale at the time of reheating [$k = k_{\rm bend} \simeq a(t_{\rm RH})H_{\rm RH}$]. The present value of bending frequency $f_{\rm bend}$ is given as

$$\begin{aligned}f_{\rm bend} &= \left(\frac{g_s(t_0)}{g_s(t_{\rm RH})}\right)^{1/3}\left(\frac{T_0}{T_{\rm RH}}\right)\frac{k_{\rm bend}}{2\pi a(t_{\rm RH})} \\ &\simeq 30\ {\rm Hz}\left(\frac{T_{\rm RH}}{10^9\ {\rm GeV}}\right).\end{aligned} \tag{19}$$

We can obtain the reheating temperature $T_{\rm RH}$ from observation of the bend in the GW spectrum at the large scale as observation of the inflationary GW background [57]. In addition, we can obtain the mass of the flat direction m_ϕ and the cutoff scale M_* from observations of energy density and peak frequency of GWs [see Eqs. (17) and (18)]. Note that other mechanisms can also generate similar GW signals (see e.g., Ref. [15]). It is unclear if we can determine these parameters only with GW experiments. Although further detailed studies are warranted for this, we leave it for future work. However, let us stress that since we are motivated by SUSY theories, collider experiments and dark matter searches can also give us complementary information of SUSY-breaking parameters.

Figure 2 shows examples of GW spectra predicted by the present mechanism. We also plot single detector sensitivities for LISA [37] and Ultimate DECIGO [38] by using the online sensitivity curve generator in Ref. [58] with the parameters in Table 7 of Ref. [59]. Note that in some parameter space we can obtain $T_{\rm RH}$ by using only one single detector for Ultimate DECIGO. One may wonder

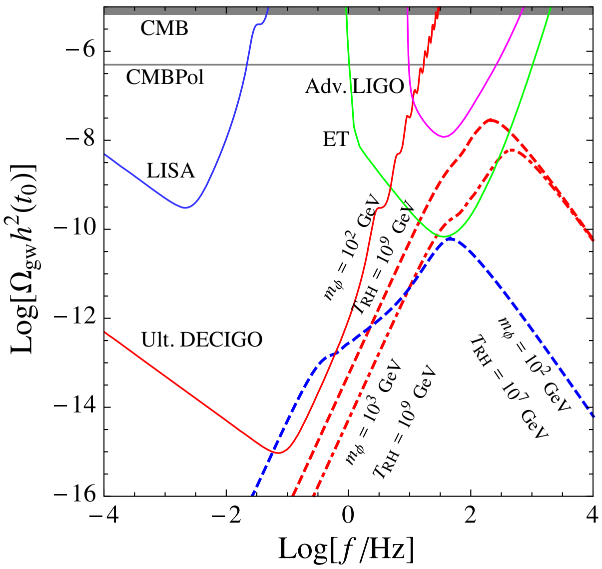

FIG. 2 (color online). GW spectra generated by cosmic strings and sensitivities of planned interferometric detectors. We plot the cases with $m_\phi = 10^2$ GeV (red dashed curve) and $m_\phi = 10^3$ GeV (red dot-dashed curve) for $T_{\rm RH} = 10^9$ GeV. We also plot the case with $m_\phi = 10^2$ and $T_{\rm RH} = 10^7$ GeV (blue dashed curve). We have assumed $M_*^2/M_{\rm Pl}^2 = 0.1$, $n = 4$, and $c'_H = a'_H = 1$.

why our sensitivity curve is different from that in Ref. [57], where they claim that Ultimate DECIGO can reach well below $\Omega_{\rm gw} \sim 10^{-16}$ for frequencies 0.1–1 Hz. This is because they assume taking two detectors correlated for a decade to obtain $T_{\rm RH}$ by observation of the inflationary GW background, while we assume a single detector. A single detector is sufficient for our purpose because GW signals from cosmic strings are much stronger than inflationary ones. We plot cross-correlation sensitivities for Advanced LIGO [39] and ET (ET-B configuration) [40], assuming two detectors are coaligned and coincident. In the figure, we take the signal-to-noise ratio SNR = 5, the angular efficiency factor $F = 2/5$, the total observation time $T = 1$ yr, and the frequency resolution $\Delta f/f = 0.1$. CMB constraints (horizontal lines) are put on the integrated energy density of GWs $\int d\log f \Omega_{\rm gw} h^2(t_0)$ [60]. We find that GW signals would be observed by ET and Ultimate DECIGO.

VI. CONCLUSIONS

We have shown that in a wide class of SUSY models, the sign of the Hubble-induced mass term for flat directions in SUSY theories generally changes just after the end of inflation, which results in the formation of cosmic strings. These cosmic strings disappear when the Hubble parameter decreases down to the soft mass of the flat direction, which implies that GWs emitted by the cosmic strings have a peak frequency related to the soft mass of the flat direction. Although other mechanisms are expected to generate similar GW signals, the observation of peak frequency can yield information of SUSY-breaking parameters in conjunction with results from collider experiments and dark matter searches. We have numerically calculated the spectrum of GWs and have shown that the resulting GW backgrounds would be measured by future GW observation experiments such as ET and Ultimate DECIGO.

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