Heterotic cosmic strings

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We show that all three conditions for the cosmological relevance of heterotic cosmic strings, the right tension, stability and a production mechanism at the end of inflation, can be met in the strongly coupled M-theory regime. Whereas cosmic strings generated from weakly coupled heterotic strings have the well-known problems posed by Witten in 1985, we show that strings arising from M5-branes wrapped around 4-cycles (divisors) of a Calabi-Yau in heterotic M-theory compactifications solve these problems in an elegant fashion.

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I. INTRODUCTION

It has been known for a long time that Cosmic Background Explorer (COBE) data require the effective or fundamental tension μ of a cosmic string to be given by $G_N\mu \simeq 10^{-6}$ if the scaling solution of the cosmic string network is assumed to be the prime source for density perturbations which seed galaxy formation. The option that cosmic strings are primarily responsible for structure formation has, however, been ruled out by more recent cosmic microwave background radiation (CMB) data. More precisely, it has been shown [1] that present CMB data [2] constrain the contribution of a cosmic string network to the CMB anisotropies to be less than 20%. This leads to a slightly tighter upper bound,

$$G_N\mu \lesssim 2 \times 10^{-7}, \tag{1.1}$$

on the cosmic string tension. The bound can equivalently be written as $\sqrt{\mu} \lesssim 5.5 \times 10^{15}$ GeV and indicates that the energy scale associated with the cosmic string tension should be roughly of the order of the grand unified theory (GUT) scale (for recent reviews on cosmic strings see [3–6]).

For the weakly coupled heterotic string, μ equals the fundamental string's tension $T = 1/2\pi\alpha'$ which is given by the string-scale squared M_s^2. Since $M_s \simeq 10^{18}$ GeV we are 2.5 orders of magnitude above the required energy scale and would hence violate the bound (1.1). Another way to see this is to remember the fact that, in the weakly coupled heterotic string, gravitational and gauge couplings are tightly related, $4\kappa_{10}^2 = \alpha' g_{10}^2$, since both originate at the level of the trilinear interactions of the closed heterotic string. This same origin also implies that both gravity and the gauge fields live in the total 10d spacetime (this no longer holds for the strongly coupled heterotic string) and therefore both couplings reduce in the same way to the corresponding 4d couplings. With $\alpha_{\mathrm{GUT}} \simeq 1/25$ being the 4d gauge coupling whose value follows from the unification of all gauge forces, we obtain

$$G_N\mu = \frac{\alpha' \alpha_{\mathrm{GUT}}}{8}\mu \simeq 8 \times 10^{-4}, \tag{1.2}$$

which clearly violates the bound (1.1). Consequently, weakly coupled heterotic fundamental strings cannot lead to viable cosmic strings, as was realized by Witten 20 years ago [7,8].

In type II theories the string scale can be lowered down to the TeV scale. This allows for a large range of cosmic string tensions below the GUT scale in compliance with the observational bound [10,11]. However, this large range for the fundamental string scale weakens the predictivity of type II cosmic strings. Their tensions might well be below observational verification. To have a more predictive framework, we will now consider the strongly coupled heterotic string where the Planck scale is fixed. The fact which makes this theory very interesting for cosmic strings is that the gravitational coupling scale which determines the M2- and M5-brane tensions,

$$\kappa_{11}^{2/9} \simeq \frac{1}{2M_{\mathrm{GUT}}}, \tag{1.3}$$

coincides roughly with the 4d GUT scale $M_{\mathrm{GUT}} \simeq 3 \times 10^{16}$ GeV [12]. Hence we can expect that the effective tensions of cosmic strings arising from suitably wrapped M2- and M5-branes might be close to the bound (1.1). This is our main reason to focus on the strongly coupled heterotic string, or heterotic M theory for short [13]. We will show in this paper that all three criteria—tension, stability, production at the end of inflation—can be satisfied in the M5-brane case.

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II. COSMIC STRING CANDIDATES FROM WRAPPED M2- AND M5-BRANES

Heterotic M theory contains only two extended objects, the M2- and the M5-branes which we are exploring as candidates for heterotic cosmic strings. The theory also contains 10-dimensional boundaries which might loosely be regarded as M9-branes. They fill, however, all of the 4-dimensional spacetime and can therefore not generate cosmic strings. For the generation of gauge cosmic strings, which we are not investigating here, this is another matter, as the Yang-Mills vector bundles are localized precisely on the M9's. It should be interesting to explore this question in the future. Generating a cosmic string from wrapped M2 or M5-branes means that these branes must extend along a timelike and a spacelike direction, t, x, into 4-dimensional spacetime.

We consider heterotic M theory compactified on $X \times \mathbf{S}^1/\mathbf{Z}_2$, where X is a Calabi-Yau threefold. The resulting flux-compactification geometry has in the simplest case a Calabi-Yau which is conformally deformed by a warp factor generated from the background $G^{(2,2,0)}$ flux [21–23] (see also [24]). We will now consider wrapping M2- and M5-branes over suitable cycles in this 7-dimensional flux-compactification background and start by listing all possible candidates for obtaining cosmic strings in 4 dimensions.

Let us begin with those configurations which are considered Bogomol'nyi-Prasad-Sommerfield state (BPS) in the flat spacetime limit. These are the M2-brane transverse to the M9's and the M5-brane parallel to them. The M2-brane which stretches along the $\mathbf{S}^1/\mathbf{Z}_2$ interval produces in the limit of vanishing orbifold length L, i.e. the weakly coupled limit, a fundamental heterotic string. Since the fundamental heterotic string is a closed string, we learn that the M2-brane world volume must have the following topology:

$$\mathrm{M}2_{\perp}: \quad \underbrace{\mathbf{R}^1 \times \mathbf{S}^1}_{\text{cosmic string loop}} \times \mathbf{S}^1/\mathbf{Z}_2, \tag{2.1}$$

giving rise to a cosmic string loop.

The parallel M5-brane needs to wrap a 4-cycle $\mathbf{\Sigma}_4$ on X to produce a stringlike object. For this we need to adopt a Calabi-Yau with nonvanishing $b_4(X) = 2h^{3,1} + h^{2,2} = h^{1,1} \neq 0$ which is the generic case. The topology of the M5-brane world volume will then be

$$\mathrm{M}5_{\parallel}: \quad \underbrace{\mathbf{R}^1 \times \mathbf{R}^1}_{\infty\text{-extended cosmic string}} \times \mathbf{\Sigma}_4, \tag{2.2}$$

where the two noncompact time and space directions are along the two M5-brane dimensions which extend into the 4-dimensional spacetime and create naturally an infinitely extended cosmic string.

One might also contemplate parallel M2-branes by wrapping the M2 not along $\mathbf{S}^1/\mathbf{Z}_2$ but instead on a 1-cycle of X. This would also create a string but can be ruled out because the Calabi-Yau threefold has vanishing first Betti number, $b_1(X) = 2h^{1,0} = 0$, hence possesses no 1-cycles on which the M2 could be wrapped (we will not consider nonsimply connected Calabi-Yau's). More interesting are the transverse M5-branes which wrap one of the $b_3(X) = 2(h^{3,0} + h^{2,1}) = 2(1 + h^{2,1}) \neq 0$ 3-cycles $\mathbf{\Sigma}_3$ and have topology

$$\mathrm{M}5_{\perp}: \quad \underbrace{\mathbf{R}^1 \times \mathbf{R}^1}_{\infty\text{-extended cosmic string}} \times \mathbf{\Sigma}_3 \times \mathbf{S}^1/\mathbf{Z}_2. \tag{2.3}$$

The resulting cosmic string would again be an infinitely extended cosmic string.

We will next derive the tensions of the cosmic string and compare them with the constraint (1.1). An important role will be played by the warped background which influences the tension. The observational bound will eliminate the $\mathrm{M}2_{\perp}$ candidate and leave us with the two M5-brane candidates.

III. COSMIC STRING TENSIONS

A. $\mathrm{M}2_{\perp}$-brane case

Let us begin with the $\mathrm{M}2_{\perp}$-brane. To determine the effective tension of the associated cosmic string, we take the Nambu-Goto part of the $\mathrm{M}2_{\perp}$-brane action

$$S_{\mathrm{M}2} = \tau_{\mathrm{M}2} \int_{\mathbf{R}^1} dt \int_{\mathbf{S}^1} dx \int_0^L dx^{11} \sqrt{-\det h_{ab}} + \ldots, \tag{3.1}$$

and integrate it over the compact dimension x^{11}. Here $a, b, \ldots = t, x, x^{11}$ and L is the length of the $\mathbf{S}^1/\mathbf{Z}_2$ interval. We adopt a static gauge for the embedding of the $\mathrm{M}2_{\perp}$ into 11-dimensional spacetime which gives us for the induced metric ($I, J = 0, \ldots, 9, 11$)

$$h_{ab} \equiv \frac{\partial X^I}{\partial x^a}\frac{\partial X^J}{\partial x^b} G_{IJ} = \delta^I_a \delta^J_b G_{IJ}. \tag{3.2}$$

The 11d metric G_{IJ} is given by the warped G-flux-compactification background sourced by the boundary M9's [21–23],

$$\begin{aligned} ds_{11}^2 &= G_{IJ} dx^I dx^J \\ &= e^{-f(x^{11})} g_{\mu\nu} dx^\mu dx^\nu + e^{f(x^{11})} (g(X)_{lm} dy^l dy^m \\ &\quad + dx^{11} dx^{11}), \end{aligned} \tag{3.3}$$

where the warp factor is given by [25]

$$e^{f(x^{11})} = (1 - x^{11} Q_v)^{2/3} \tag{3.4}$$

with visible $M9$-brane charge

$$Q_v = -\frac{1}{8\pi V_v}\left(\frac{\kappa_{11}}{4\pi}\right)^{2/3} \int_{X_v} J \wedge \left(\mathrm{tr} F \wedge F - \frac{1}{2}\mathrm{tr} R \wedge R\right) \tag{3.5}$$

which sources the $G^{(2,2,0)}$ flux component. X_v and V_v denote the Calabi-Yau and its volume at the location of the visible M9, J its Kähler-form and F resp. R the Yang-Mills and curvature 2-forms, again on the visible M9.

Notice that we are taking the flux background which incorporates only the backreaction of the M9's but not that of extra $M5_2$-branes in the bulk. The extra $M5_2$-branes would wrap genus zero holomorphic 2-cycles on X and fill all of 4-dimensional spacetime so should not be confused with the $M5_{\parallel}$-, $M5_{\perp}$-brane candidates for cosmic strings. Though the backreaction of the $M5_2$-branes is known [21–23], their neglect is justified when we want to focus on a cosmological epoch at the end of inflation or even later which is the time when the cosmic strings are produced and observed. In the proposal for heterotic M-theory inflation made in [26,27], which we will use here, the inflationary dynamics relies on the interactions between several $M5_2$-branes in the bulk. Towards the end of inflation, the 11-dimensional bulk, however, gets cleared of its $M5_2$-branes which coalesce with the boundary M9's. This justifies the neglect of the $M5_2$-branes in the flux background. Let us also note that the addition of $M5_2$-branes would weaken the tight relation between the GUT and gravity sector which relates so successfully the standard values for $M_{\rm GUT} \simeq 3\times 10^{16}$ GeV and $\alpha_{\rm GUT} \simeq 1/25$ to the observed value for Newton's constant G_N.

We can now explicitly integrate over x^{11} with the result that the M2-brane action becomes the cosmic string action,

$$S_{\rm M2} = \mu_{\rm M2}\int_{\mathbf{R}^1} dt \int_{\mathbf{S}^1} dx\sqrt{-g_{tt}g_{xx}} + \cdots, \tag{3.6}$$

with tension determined by the warp factor and length L of the $\mathbf{S}^1/\mathbf{Z}_2$ interval,

$$\mu_{\rm M2} = \tau_{\rm M2}\int_0^L dx^{11} e^{-f(x^{11})/2} = \frac{3\tau_{\rm M2}}{2Q_v}(1-(1-LQ_v)^{2/3}). \tag{3.7}$$

To evaluate the value, let us remember that the correct value of the 4d Newton's constant requires L to be of critical length L_c which is given in terms of the M9 charge by [21,22]

$$L_c \equiv 1/Q_v. \tag{3.8}$$

We should therefore use $L \simeq L_c$ for the evaluation of the cosmic string's tension. To evaluate the tension, let us express all quantities in terms of the 11-dimensional gravitational coupling constant κ_{11}. Based on phenomenological reasoning, the critical length will be given by [12,22]

$$L_c \simeq 12\kappa_{11}^{2/9}. \tag{3.9}$$

With the M2-brane tension $\tau_{\rm M2} = M_{11}^3/(2\pi)^2$, and the defining relation $2\kappa_{11}^2 = (2\pi)^8/M_{11}^9$ for the 11d Planck mass M_{11}, we obtain for the string's tension

$$\mu_{\rm M2} = \frac{3\tau_{\rm M2}}{2Q_v} = 3L_c\left(\frac{\pi}{2\kappa_{11}}\right)^{2/3} \simeq 9(2^{10}\pi^2)^{1/3}M_{\rm GUT}^2. \tag{3.10}$$

For the last expression we have used the relations (1.3) and (3.9). Since $\mu_{\rm M2}^{1/2}$ turns out to be larger than the GUT scale, it is clear that the string's tension comes out too large. This becomes evident when we finally evaluate

$$G_N\mu_{\rm M2} \simeq 1.2\times 10^{-3} \tag{3.11}$$

with $M_{\rm GUT} \simeq 3\times 10^{16}$ GeV which is in clear conflict with the observational bound (1.1). Also, considering a slightly smaller length $L = 11\kappa_{11}^{2/9} = L_c - \kappa_{11}^{2/9}$, which could still be stabilized at the end of inflation, would only decrease the tension by a factor of 0.8 which is not enough. The $M2_{\perp}$ candidates are therefore ruled out as viable cosmic strings.

B. $M5_{\parallel}$-brane case

Let us now turn to the $M5_{\parallel}$ cosmic strings. The Nambu-Goto term of the $M5_{\parallel}$-brane action reads

$$S_{M5_{\parallel}} = \tau_{\rm M5}\int_{\mathbf{R}^1} dt \int_{\mathbf{R}^1} dx \int_{\mathbf{\Sigma}_4} d^4y\sqrt{-\det h_{ab}} \tag{3.12}$$

where $a, b, \ldots = t, x, y^1, y^2, y^3, y^4$. Adopting again static gauge for its embedding, we have to integrate over the 4-cycle $\mathbf{\Sigma}_4$ to obtain the action for the cosmic string,

$$S_{M5_{\parallel}} = \mu_{M5,\parallel}\int_{\mathbf{R}^1} dt \int_{\mathbf{R}^1} dx\sqrt{-g_{tt}g_{xx}} \tag{3.13}$$

with the string tension given by

$$\mu_{M5_{\parallel}} = \tau_{\rm M5}e^{f(x_{\rm M5}^{11})}\int_{\mathbf{\Sigma}_4} d^4y\left(\prod_{i=1,\ldots,4} g(X)y^iy^i\right)^{1/2} = \tau_{\rm M5}\left(1-\frac{x_{\rm M5}^{11}}{L_c}\right)^{2/3} V_{\mathbf{\Sigma}_4}. \tag{3.14}$$

Here $0 \le x_{\rm M5}^{11} \le L$ denotes the position of the $M5_{\parallel}$ along the $\mathbf{S}^1/\mathbf{Z}_2$ orbifold. It will be convenient to write the volume of the 4-cycle $V_{\mathbf{\Sigma}_4}$ in terms of a dimensionless radius $r_{\mathbf{\Sigma}_4}$ by rescaling with the radius R_v of X on the visible boundary, i.e. the undeformed initial Calabi-Yau radius

$$V_{\mathbf{\Sigma}_4} = (r_{\mathbf{\Sigma}_4}R_v)^4. \tag{3.15}$$

Typically, one would expect for a more or less isotropic Calabi-Yau that $r_{\mathbf{\Sigma}_4} \lesssim 1$. For highly anisotropic compactification spaces, it could be larger.

To evaluate the tension's value, we need to employ another standard relation [12,22],

$$R_v \equiv V_v^{1/6} = 1/M_{\rm GUT}. \tag{3.16}$$

Using this, the definition of the $M5_{\parallel}$-brane's tension, $\tau_{M5} = M_{11}^6/(2\pi)^5$, plus (1.3), we arrive at

$$\mu_{\mathrm{M5}_{\|}} = 64\left(\frac{\pi}{2}\right)^{1/3}\left(1-\frac{x^{11}_{\mathrm{M5}}}{L_c}\right)^{2/3} M^2_{\mathrm{GUT}} r^4_{\mathbf{\Sigma}_4}. \quad (3.17)$$

Numerically this leads to the following result:

$$G_N\mu_{\mathrm{M5}_{\|}} = 4.7\times 10^{-4}\left(1-\frac{x^{11}_{\mathrm{M5}}}{L_c}\right)^{2/3} r^4_{\mathbf{\Sigma}_4}. \quad (3.18)$$

We will subsequently see that the production of the $\mathrm{M5}_{\|}$ cosmic strings will happen towards the end of inflation essentially on the hidden M9 when L gets stabilized near L_c [30,31]. Taking therefore, say, $x^{11}_{\mathrm{M5}} = L \simeq 11\kappa_{11}^{2/9} = L_c - \kappa_{11}^{2/9}$, we obtain $G_N\mu_{\mathrm{M5}_{\|}} = 8.9\times 10^{-5} r^4_{\mathbf{\Sigma}_4}$. A radius $r_{\mathbf{\Sigma}_4} \le 0.22$ would then already be enough to satisfy the observational constraint. Hence the $\mathrm{M5}_{\|}$ easily passes the tension constraint. The positioning of the $\mathrm{M5}_{\|}$-brane on the hidden boundary is also supported by the fact that M5-branes can only wrap 4-cycles which carry no G flux [32]. In general, this is the case on either the visible or hidden M9 boundary where the $G^{(2,2,0)}$ flux vanishes as a direct consequence of the $\mathbf{Z}_2$ symmetry of the background.

C. $\mathrm{M5}_\perp$-brane case

Let us finally come to the $\mathrm{M5}_\perp$ cosmic strings. We start from the $\mathrm{M5}_\perp$-brane action

$$S_{\mathrm{M5}_\perp} = \tau_{\mathrm{M5}}\int_{\mathbf{R}^1} dt\int_{\mathbf{R}^1} dx\int_0^L dx^{11}\int_{\mathbf{\Sigma}_3(x^{11})} d^3y\sqrt{-\det h_{ab}}. \quad (3.19)$$

Integrating over the compact dimensions gives the cosmic string action

$$S_{\mathrm{M5}_\perp} = \mu_{\mathrm{M5}_\perp}\int_{\mathbf{R}^1} dt\int_{\mathbf{R}^1} dx\sqrt{-g_{tt}g_{xx}} \quad (3.20)$$

with string tension

$$\mu_{\mathrm{M5}_\perp} = \tfrac{3}{5}\tau_{\mathrm{M5}}(1-(1-L/L_c)^{5/3})L_c V_{\mathbf{\Sigma}_3}. \quad (3.21)$$

Again it will be convenient to express the volume of the 3-cycle $V_{\mathbf{\Sigma}_3}$ through a dimensionless radius $r_{\mathbf{\Sigma}_3}$ defined by

$$V_{\mathbf{\Sigma}_3} = (r_{\mathbf{\Sigma}_3}R_v)^3. \quad (3.22)$$

With the standard relations used earlier, we arrive then at

$$\mu_{\mathrm{M5}_\perp} = \frac{72}{5}\left(\frac{\pi}{2}\right)^{1/3}(1-(1-L/L_c)^{5/3})M^2_{\mathrm{GUT}} r^3_{\mathbf{\Sigma}_3}, \quad (3.23)$$

which gives the result

$$G_N\mu_{\mathrm{M5}_\perp} = 1.1\times 10^{-4}(1-(1-L/L_c)^{5/3}) r^3_{\mathbf{\Sigma}_3}. \quad (3.24)$$

Again, for a value $L = 11\kappa_{11}^{2/9}$, we obtain $G_N\mu_{\mathrm{M5}_\perp} = 1.1\times 10^{-4} r^3_{\mathbf{\Sigma}_3}$. Hence the observational constraint can be satisfied for $r_{\mathbf{\Sigma}_3} \le 0.12$. This still seems a rather mild constraint on the average radius of the 3-cycle $\mathbf{\Sigma}_3$. We can therefore conclude that the $\mathrm{M5}_\perp$ cosmic strings also pass the tension test. We will next analyze the stability of our two M5-brane candidates.

IV. STABILITY

A. Classical stability

Cosmic strings resulting from fundamental heterotic strings were found in [7] to be unstable. The reason was that these cosmic strings are axionic strings with $\mathbf{S}^1$ topology which bound domain walls. Because of the domain wall tension which is proportional to the area they span, these axionic strings will quickly shrink. Hence they cannot become macroscopically large.

We will, at first sight, encounter the same instability for cosmic strings resulting from wrapped M2- or M5-branes in heterotic M theory. This is because these branes are charged under the 3-form C_3 resp. dual 6-form potential C_6, which, when reduced over the appropriate cycle that the brane wraps, becomes a 2-form potential $C_{[2]}$ in 4 dimensions. Since the dual of this 2-form gives an axion ϕ via

$$dC_{[2]} = \star_4 d\phi, \quad (4.1)$$

it seems that cosmic strings created by wrapping M2- or M5-branes cannot grow to cosmic size due to their coupling to the axion ϕ. To avoid this conclusion one needs to remove the massless axion. We will see that this will only be possible for the $\mathrm{M5}_{\|}$ cosmic string candidate and requires it to be on the hidden M9. Hence the $\mathrm{M5}_\perp$ cosmic string candidate will be ruled out, as it suffers from the domain wall instability and therefore quickly shrinks to microscopic size. Let us now explain how and under which conditions the massless axion gets removed.

For this, let us remember first that the presence of the boundaries in heterotic M theory leads to a modification of its 4-form field strength G on the boundaries [33]. This modification involves the Yang-Mills and Lorentz Chern-Simons 3-forms ω_Y, ω_L, and one finds on the hidden boundary [34] at $x^{11} = L$

$$G_4 = dC_3 + c\kappa_{11}^{2/3}\left(\omega_Y - \frac{1}{2}\omega_L\right)\delta(x^{11}-L)\wedge dx^{11},$$
$$c = \frac{\sqrt{2}}{(4\pi)^{5/3}}. \quad (4.2)$$

To avoid carrying around the delta function, let us write this in 10d notation in terms of the Neveu-Schwarz 3-form field strength H on the hidden boundary (where $H_{ABC} = G_{11ABC}$, $B_{AB} = C_{11AB}$)

$$H_3 = dB_2 - \frac{c\kappa_{11}^{2/3}}{2L}\left(\omega_Y - \frac{1}{2}\omega_L\right). \quad (4.3)$$

Since $\alpha' = 2c\kappa_{11}^{2/3}/L$ [12], we recognize the familiar α' correction of the weakly coupled heterotic string, with the difference of the factor $1/2$ which arises from the separation of the boundaries. Plugging this field strength into the hidden boundary $\mathcal{M}_h^{10}$ kinetic term

$$-\frac{L}{2\kappa_{11}^2}\int_{\mathcal{M}_h^{10}} H_3 \wedge \star_{10} H_3 \tag{4.4}$$

leads upon dualization $dC_6 = \star_{10} dB_2$ to the coupling

$$\frac{c}{2\kappa_{11}^{4/3}}\int_{\mathcal{M}_h^{10}} C_6 \wedge \left(\mathrm{tr} F\wedge F - \frac{1}{2}\,\mathrm{tr} R\wedge R\right). \tag{4.5}$$

We know that, in order to stabilize the hidden boundary close to the phenomenologically relevant length L_c after inflation, the hidden E_8 gauge symmetry must be broken to a gauge group of smaller rank [31]. This will typically provide us with some $U(1)$ gauge symmetries on the hidden M9. Let us pick one of these and denote its field strength $\mathcal{F}_2 = dA_1$. Moreover, let us assume a nonvanishing gauge flux $\int_{C_2} F \neq 0$ over some 2-cycle on X. Let us consider the coupling term together with the kinetic terms in the 11-dimensional action

$$-\frac{1}{2\times 7!\kappa_{11}^2}\int_{\mathcal{M}^{11}} |dC_6|^2 + \frac{c}{2\kappa_{11}^{4/3}}\int_{\mathcal{M}_h^{10}} C_6 \wedge \left(\mathrm{tr} F\wedge F - \frac{1}{2}\,\mathrm{tr} R\wedge R\right) - \frac{1}{4g_{10}^2}\int_{\mathcal{M}_h^{10}} |F|^2. \tag{4.6}$$

Here the 10-dimensional gauge coupling g_{10} is fixed in terms of the gravitational coupling as $g_{10}^2 = (2^7\pi^5)^{1/3}\kappa_{11}^{4/3}$ [33]. After a reduction to four dimensions these terms will give a contribution (we will not consider the curvature term $\mathrm{tr}R\wedge R$ further)

$$-\frac{1}{2}\int_{\mathcal{M}^4} |dC_{[2]}|^2 + m\int_{\mathcal{M}^4} C_{[2]}\wedge \mathcal{F}_2 - \frac{1}{2}\int_{\mathcal{M}^4} |\mathcal{F}_2|^2 \tag{4.7}$$

to the 4-dimensional action. The mass parameter m is given by

$$m = \frac{(7!)^{1/2}}{2^{8/3}\pi^{5/6}} \times \frac{\kappa_{11}^{1/3} L_{\mathrm{top}}^4}{(L\langle V\rangle V_h)^{1/2}}, \tag{4.8}$$

where $\langle V\rangle$ denotes the Calabi-Yau volume averaged over the $\mathbf{S}^1/\mathbf{Z}_2$ interval, V_h represents the Calabi-Yau volume at the location of the hidden boundary, and the length L_{top} will be defined next. To arrive at this expression, we have set

$$\int_{\mathcal{M}_h^{10}} C_6 \wedge \mathrm{tr}(\mathcal{F}_2\wedge F) = L_{\mathrm{top}}^4 \int_{\mathcal{M}^4} C_{[2]}\wedge \mathcal{F}_2 \tag{4.9}$$

and then rescaled

$$\mathcal{F}_2 \to \left(\frac{V_h}{2g_{10}^2}\right)^{1/2}\mathcal{F}_2, \tag{4.10}$$

$$C_{[2]} \to \left(\frac{2\langle V\rangle L}{7!\kappa_{11}^2}\right)^{1/2} C_{[2]} \tag{4.11}$$

such that the 4-dimensional fields $C_{[2]}$, A_1 receive a canonical mass dimension one. The volume and length factors which enter the rescaling originate from the ordinary reduction of the metric-dependent kinetic terms for C_6 and A_1 from 11 resp. 10 to 4 dimensions. The length parameter L_{top} which stems from the reduction of the metric-independent topological coupling term characterizes the localization of the gauge flux F and C_6 on X.

It is now straightforward to demonstrate [35] that this action implies the absence of the axion ϕ which we will show next. The field equations for A_1 and $C_{[2]}$ which result from the action (4.7) are

$$d\star_4 dA_1 = -m dC_{[2]}, \tag{4.12}$$

$$d\star_4 dC_{[2]} = -m\mathcal{F}_2. \tag{4.13}$$

We can solve the second equation by

$$dC_{[2]} = \star_4(d\phi - mA_1), \tag{4.14}$$

which defines the dual axion field ϕ. Plugging this solution back into the field equation for A_1 gives

$$d\star_4 dA_1 = \star_4(-m d\phi + m^2 A_1). \tag{4.15}$$

For the ground state in which $\phi = 0$ or by picking a gauge which sets $d\phi = 0$, this result shows that A_1 has acquired a mass m. Alternatively, one might plug the solution back into the action (4.7). Then the coupling term gives us a mass term for A_1

$$m\int_{\mathcal{M}^4} C_{[2]}\wedge dA_1 = \int_{\mathcal{M}^4}(mA_1\wedge \star_4 d\phi - m^2 A_1\wedge \star_4 A_1). \tag{4.16}$$

Furthermore, we infer from (4.14) that ϕ must transform nonlinearly under A_1 gauge transformations

$$\delta A_1 = d\Lambda, \qquad \delta\phi = -m\Lambda. \tag{4.17}$$

The proper interpretation of these results is that the $U(1)$ gauge field swallows the axion ϕ, gains a further degree of freedom, and becomes massive, i.e. $A_1 \to A_1 - d\phi/m$. Since the axion gets removed in this Higgsing, there is no domain wall anymore which would prevent the cosmic string from growing. Let us note that m grows when the hidden boundary comes close to the critical length L_c where V_h would classically vanish and is expected to quantum mechanically reach Planck size [38] $l_{11}^6 \simeq (\kappa_{11}^{2/9}/5)^6$. Since towards the end of the inflationary mechanism of [26] the hidden boundary gets indeed stabilized close to L_c, where V_h becomes small, through the stabilization mechanisms developed in [30,31] we notice that the removal of the axion domain wall will be particularly effective towards the end of inflation when m becomes large.

For which of our cosmic string candidates, $M5_{\|}$, $M5_{\perp}$, does this stabilization mechanism apply? The gauge fields F are localized on the boundary and therefore the initial coupling (4.5) will only be nonvanishing for a parallel $M5_{\|}$-brane which moreover has to be localized on the hidden boundary. The transverse $M5_{\perp}$ which stretches orthogonal to $\mathcal{M}_h^{10}$ along $\mathbf{S}^1/\mathbf{Z}_2$ cannot have this coupling. It will therefore maintain its domain wall instability and will consequently quickly shrink to microscopic size. This might have been anticipated because the $M5_{\perp}$ is a non-BPS object in flat 11-dimensional spacetime. We are therefore left with a unique cosmic string candidate, a parallel $M5_{\|}$-brane on the hidden boundary.

Let us now come to a second potential instability which is the breaking of the $M5_{\|}$ cosmic string on the hidden boundary. Since the end points which are produced when the string breaks are still connected by flux lines, one can think of this breaking as the $M5_{\|}$-brane dissolving in the M9. One has to compare the gauge flux $\int_{\mathcal{C}_2} F$ which is transverse to the $M5_{\|}$-brane with the kinetic energy density $\int_X F \wedge \star_6 F$ on X. By counting dimensions, one would conclude that it might be energetically favorable for the flux to expand along X and therefore the cosmic string might break.

The reason why this conclusion should not hold is very simple. Notice that the argument so far implicitly assumed that X is large enough in order to provide space for the flux to spread along X. This, however, is not the case precisely on the hidden M9. As we will review later, L, and therefore the hidden M9, gets stabilized towards the end of inflation close to L_c. The characteristic feature of L_c is that it is the length at which the volume of X shrinks classically to a point. Therefore the flux has no space to spread along X when the $M5_{\|}$-brane is on (or close to) the hidden M9. Another argument against the breaking of the string, even at finite size X volumes, might also come from the nice solution of the breaking instability for a D1 on a D3-brane presented in [37]. Here, as well as in our case, we have a flux $\int_{\mathcal{C}_2} F \neq 0$ transverse to the cosmic resp. D1-string. Since we do not have, however, a sizable volume for X, we will not explore this possibility further here.

So it remains to analyze whether there can be breakage of the $M5_{\|}$ cosmic string in the four noncompact directions. Here, let us note that the $M5_{\|}$ cosmic strings, when located on the hidden M9, lead in 4 dimensions to an effective Abelian Higgs model whose $U(1)$ is Higgsed. Consequently, Abrikosov-Nielsen-Olesen type flux tubes [39] will form which carry magnetic flux of the Higgsed $U(1)$. These flux tubes, in which the field strength falls off exponentially with radial distance, cannot decay because they are topologically stable. It is these flux tubes which represent the $M5_{\|}$ cosmic strings in the effective 4-dimensional theory and show that they are also stable with respect to breakage along the noncompact directions. One might worry that, at high energies when the gauge theory on the hidden M9 is expected to restore a GUT symmetry [40] with a corresponding embedding of the $U(1)$ into the unified gauge group, the flux tubes might break. The reason is that GUT theories possess monopoles such that the flux tube can start on a monopole and end on an antimonopole, thus making it unstable against monopole pair production. An estimate of the monopole pair creation rate via the Schwinger pair production calculation shows, however, that this rate is suppressed by a factor $\exp[-(\pi M^2/\mu_{M5_{\|}})]$ with M being the monopole mass. We expect the $M5_{\|}$ cosmic string's tension $\mu_{M5_{\|}}$ to be far smaller than the monopole's mass, again due to its warp-factor suppression. Therefore, the scale of the monopole mass should easily be an order of magnitude larger than the scale of the string's tension which is enough to render the flux tubes effectively stable on cosmological time scales [4]. Before describing how the parallel $M5_{\|}$-branes are produced when inflation comes to an end, we will now briefly address the stability of M2-branes and quantum instabilities.

Though we have seen that the tension of an M2 cosmic string violates the observational bound and M2 cosmic strings are consequently ruled out, let us nevertheless include the stability discussion for a hypothetical M2 cosmic string. In this case there is a similar coupling, the well-known [33]

$$\frac{\sqrt{2}}{(4\pi)^3(4\pi\kappa_{11}^2)^{1/3}} \int_{\mathcal{M}^{11}} C_3 \wedge X_8(F,R) \tag{4.18}$$

where

$$X_8(F,R) = -\frac{1}{4}\Big(\mathrm{tr}F^2 - \frac{1}{2}\,\mathrm{tr}R^2\Big)^2 + \Big(-\frac{1}{8}\,\mathrm{tr}R^4 + \frac{1}{32}(\mathrm{tr}R^2)^2\Big). \tag{4.19}$$

Combining it with the kinetic terms for C_3 and F can, once again, generate the desired effective 4d coupling $\int_{\mathcal{M}^4} C_{[2]} \wedge \mathcal{F}_2$. This time it requires an orthogonal $M2_{\perp}$-brane because F is localized on the boundary M9's. Assuming a nonzero higher instanton charge $\int_X (F \wedge F \wedge F) \neq 0$ on X, we would likewise remove the axion and the associated domain wall through Higgsing of the 4-dimensional $U(1)$. This time we have a topological charge $\int_X (F \wedge F \wedge F)$ on X which we need to compare to the energy density term $\int_X F \wedge \star_6 F$. Counting dimensions, we would conclude that it is energetically favorable for the flux to shrink. Hence, the hypothetical M2 cosmic string would not break up, as it cannot transform into flux which can spread out over the M9. We will later also see that transverse branes will not be produced at the end of inflation. The stability of the $M2_{\perp}$-brane will therefore not imply its presence.

B. Quantum stability

One might ask whether the $M5_{\|}$ cosmic strings could decay quantum mechanically via some nonperturbative effect. With only M2 and M5-brane instantons available, this would require that either of them must be able to couple to the $M5_{\|}$-brane. For the M2 instantons [41] to mediate a force, they would need to wrap a genus zero holomorphic 2-cycle $\mathbf{\Sigma}_2^0$ on the divisor $\mathbf{\Sigma}_4$. Hence, if the divisor $\mathbf{\Sigma}_4$ does not contain any such 2-cycles $\mathbf{\Sigma}_2^0$, the $M5_{\|}$-brane and thus the cosmic string would not feel a force mediated by M2 instantons. Moreover, no M5 instantons i.e. M5-branes which wrap the complete X at some fixed location along the $\mathbf{S}^1/\mathbf{Z}_2$ can attach to the $M5_{\|}$-branes because the M5 instantons would need two more compact dimensions than the divisor which the $M5_{\|}$ wraps can provide. Consequently, M5 instantons will not be able to exert a force on the $M5_{\|}$-branes. Therefore, with respect to M2 or M5 instanton decay, the $M5_{\|}$ cosmic strings are stable as long as the divisor $\mathbf{\Sigma}_4$ does not contain any genus zero holomorphic 2-cycles $\mathbf{\Sigma}_2^0$.

C. Relation to other types of cosmic strings

Cosmic D-strings which arise from the tachyon condensation of a brane-antibrane Dp-$\bar{\text{D}}$p pair have *a priori* a very different fundamental description from the heterotic cosmic strings originating from wrapped M5-branes. At the level of the effective 4-dimensional description there are, however, striking similarities. Let us consider for definiteness a D3-$\bar{\text{D}}$3 pair on whose world volume a D1-string forms as a tachyonic vortex [43]. The tachyon in the open string spectrum of the D3-$\bar{\text{D}}$3 system is charged under the diagonal combination of the two $U(1)$'s. When the tachyon condenses in a topologically nontrivial vacuum, the diagonal $U(1)$ is Higgsed. The effective picture [44] of the created D1-string is a topologically stable vortex solution which carries magnetic flux of the Higgsed $U(1)$ similar to an Abrikosov-Nielsen-Olesen flux tube [39]. The Ramond-Ramond charge of the D1-string stems from a Wess-Zumino coupling

$$\int_{\mathrm{D3}\text{-}\bar{\mathrm{D}}3} \mathcal{F}_2 \wedge C_2 \tag{4.20}$$

on the D3-$\bar{\text{D}}$3 world volume. Here, $\mathcal{F}_2$ denotes the field strength of the diagonal $U(1)$ and C_2 the Ramond-Ramond 2-form. In 4 dimensions the D1-string represents a cosmic string [36]. Hence, together with the kinetic terms for the gauge potential and C_2, we arrive at an effective action which is formally the same as in (4.7). Consequently, both the heterotic cosmic strings and the type II cosmic D-strings have the same effective description in terms of Abrikosov-Nielsen-Olesen type flux tubes. Indeed the analogy between both can be extended further as we will now indicate.

Solitonic descriptions of cosmic superstrings have been given in [45,46] for heterotic string motivated models and in [44,47–49] for D-strings. Although the low-energy effective actions are very similar in both cases, they differ by a dilaton-independent D-term contribution from a Fayet-Iliopoulos term ξ of the Higgsed $U(1)$. This Fayet-Iliopoulos term ξ was not obvious and therefore was omitted in the heterotic models [45,46] while it was included for the type II D1-string, being proportional to the D3-brane tension [44]. The presence of this term is crucial as it allows us to construct solitonic supersymmetric solutions free of singularities [44]. With the construction of heterotic cosmic strings in terms of wrapped M5-branes, it is natural to guess that the $M5_{\|}$ tension could provide this Fayet-Iliopoulos term on the heterotic side. Furthermore, one might wonder whether the effective heterotic M-theory action (4.7) could be extended to include a tachyon like in the effective D3-$\bar{\text{D}}$3 or D1-D3 descriptions with the tachyon playing the role of the Higgs field. This seems to be indeed the case. Similar to the type II D3-$\bar{\text{D}}$3 or D1-D3 systems where the tachyon appears when both branes are close to each other, there are fields Φ in heterotic M theory coming from M2-branes stretching between the $M5_{\|}$-brane and the hidden M9. These fields acquire a negative mass squared and hence indeed become tachyonic when the $M5_{\|}$-brane comes close to the M9 [28].

It might also be interesting to study whether viable cosmic strings originating from wrapped M5-branes may also arise in M-theory compactifications on G_2 manifolds. We will mention just a few aspects and leave a full investigation to future work. First, in contrast to the heterotic M-theory case, G_2 compactifications preserving an $N = 1$ supersymmetry must have zero G flux and hence possess no warping [50,51]. The smallness of the cosmic string tension must therefore arise from a combination of a low (as compared to the 4-dimensional Planck scale) fundamental scale $1/\kappa_{11}^{2/9}$ together with the presence of a 4-cycle of sufficiently small volume. Indeed for special cases [52] a low fundamental scale $1/\kappa_{11}^{2/9}$ close to the GUT scale has been confirmed. Second, phenomenologically viable G_2 compactifications with non-Abelian gauge groups of type A, D, or E and charged chiral matter require the presence of a 3-dimensional locus Q of A, D, or E orbifold singularities on the G_2 manifold. Q itself is smooth but the normal directions to Q have a singularity. It remains, however, an open problem [53] to construct compact G_2 manifolds with such singularities. Consequently, the full effective 4-dimensional theory is not known to date. Anomaly considerations [54] reveal, in the case of $A_n = SU(n + 1)$ gauge groups, a 7-dimensional interaction term

$$\int_{\mathcal{M}^4\times Q} K \wedge \Omega_5(A) \tag{4.21}$$

with K the 2-form field strength of a $U(1)$ gauge field which is part of the normal bundle to Q, and $\Omega_5(A)$ the

Chern-Simons 5-form satisfying $d\Omega_5(A) = \mathrm{tr} F \wedge F \wedge F$. This term does not lead, in contrast to the heterotic M-theory case with Green-Schwarz anomaly cancelling terms, to a coupling of type (4.20) needed to gauge away the axion and therefore the domain wall instability of the M5-brane cosmic string. The stability of M5-brane wrapped cosmic strings is therefore not clear in M theory on G_2 manifolds. One should also add that a viable model of inflation arising from such M-theory compactifications has yet to be constructed.

V. END OF INFLATION

So far we have systematically analyzed which cosmic string candidates pass the tension constraint and the stability criterion. The only candidate left over is a parallel $\mathrm{M5}_\parallel$-brane localized on the hidden boundary. It remains to clarify whether these branes can also be produced towards the end of inflation. Let us therefore now briefly provide some background on the end of heterotic M-theory inflation following [26].

The inflationary phase is driven through nonperturbative interactions between several $\mathrm{M5}_2$-branes distributed along the $\mathbf{S}^1/\mathbf{Z}_2$ interval. Initially close together, the repulsive interactions between neighboring $\mathrm{M5}_2$-branes drag them towards the boundaries. This characterizes the inflationary phase. The fact that many $\mathrm{M5}_2$-branes are present enhances the Hubble friction and leads to an M-theory realization of the assisted inflation idea [55] with parametrically small slow-roll parameters. As long as the distance between the $\mathrm{M5}_2$-branes stays smaller than the orbifold length L, the resulting potential assumes the required simple exponential form [56].

This changes at the end of inflation. Here the distances between the $\mathrm{M5}_2$-branes have grown to a size comparable to that of the $\mathbf{S}^1/\mathbf{Z}_2$ length L itself and further contributions to the dynamics of the system become equally relevant. These contributions are as follows: a repulsive open M2 instanton force mediated between the two boundary M9's, gaugino condensation, and fluxes. Let us detail this a bit more. The fact that at this stage the repulsive M9-M9 interaction becomes noticeable causes L to grow. Characteristic for heterotic M theory, a growing L implies a growing gauge coupling on the hidden M9. This is a consequence of the theory's warped flux-compactification background [21–24]. Hence, towards the end of inflation the hidden gauge theory becomes strongly coupled, which triggers gaugino condensation. As a consequence of gaugino condensation, a nonvanishing Neveu-Schwarz H flux will be induced on the hidden M9. This is due to a specific perfect square structure within the heterotic action which combines gaugino condensation and H flux [57] (recent discussions can also be found in [58–61]).

The great importance of these additional contributions to the potential which enter the stage only at the end of inflation—M9-M9 interaction, gaugino condensation, and H flux—lies in the fact that they will stabilize the $\mathbf{S}^1/\mathbf{Z}_2$ length ("dilaton") and the Calabi-Yau volume (see e.g. [58–66]). Most relevant for us will be the $\mathbf{S}^1/\mathbf{Z}_2$ length L. Furthermore, in vacua with positive vacuum energy, L will be stabilized close to its critical length L_c which is the length at which the hidden Calabi-Yau volume vanishes classically. This can be achieved either with the help of one remaining position-stabilized $\mathrm{M5}_2$-brane in the bulk [30] or by breaking the E_8 gauge symmetry on the hidden boundary [31,67]. A stabilization close to L_c is actually necessary to obtain a realistic value for Newton's constant and a supersymmetry breaking scale close to the TeV scale [31]. The stabilization of L close to $L_c = 12\kappa_{11}^{2/9}$, say in a regime

$$L_c - \kappa_{11}^{2/9} \le L \le L_c, \tag{5.1}$$

has, however, an immediate impact on the cosmic string tensions derived earlier. Let us focus on the viable $\mathrm{M5}_\parallel$ cosmic string where $x^{11}_{\mathrm{M5}} = L$ because we have seen that only on the boundary [69] can it be freed of its domain wall instability. From (3.18) we see that for $L \to L_c$ these cosmic strings can become nearly tensionless. Such a low tension is only possible through the warp factor of the background which contributes the $(1 - x^{11}_{\mathrm{M5}}/L_c)^{2/3}$ suppression factor to (3.18).

Let us conclude this section by stressing the salient feature of this quick review. Namely, we can influence the tension of the $\mathrm{M5}_\parallel$ cosmic string by the value at which L will be stabilized at the end of inflation. Realistic stabilizations require a stabilization close to L_c which lowers the cosmic string's tension considerably.

VI. PRODUCTION

Earlier we found that $\mathrm{M2}_\perp$ cosmic strings would violate the observational bound on the cosmic string's tension. It is therefore satisfying to see that they are not being produced when inflation comes to an end. This is due to the fact that their production would exceed the energy threshold available at this time which certainly lies below M_{GUT}. We had further seen that the tension of $\mathrm{M5}_\parallel$-branes is small enough so that they can reach cosmic size once they are produced. In this section we will qualitatively describe a mechanism which leads to the production of these heterotic cosmic strings.

The model of inflation of [70] is based on the dynamics of a pair of D3- and anti-D3-branes. Towards the end of inflation the distance between the brane and the antibrane goes to zero resulting in their annihilation. It has been argued in [36] that this annihilation results in the creation of D1-branes which can reach a cosmic size.

The mechanism leading to cosmic string production in our scenario is rather different and is based on the strongly time-dependent background which originates at the end of the inflationary process [71–74]. The heterotic M-theory inflation model presented in [26] is based on the dynamics of a set of $M5_2$-branes which, towards the end of inflation, approach the boundaries of the $\mathbf{S}^1/\mathbf{Z}_2$ interval. When the $M5_2$-branes hit the boundaries, the background becomes strongly time dependent and at this point the inflaton field starts performing rapid coherent oscillations with a Planck sized amplitude. Precisely, these oscillations provide the source of energy to pair produce strings of low tension. The production rate for these strings was evaluated in [71,72] from the physical state constraint for the string states

$$L_0|\text{physical}\rangle = 0, \tag{6.1}$$

which was rewritten as a differential equation for a string state $\chi(t)$ whose oscillation frequency $\omega(t)$ is sourced by the inflaton

$$\ddot{\chi} + \omega(t)\chi = 0. \tag{6.2}$$

It turns out that the pair produced strings cannot be fundamental strings as their tension would be of the order of the 4-dimensional Planck scale, roughly $M_{\text{Pl}} \simeq 10^{18}$ GeV, several orders of magnitude above the typical inflaton mass $m_{\text{inf}} \simeq 10^{13}$ GeV.

Nonperturbative strings would be the alternative and these are precisely the objects produced at the end of our inflationary process. Indeed, our candidates for cosmic strings are not fundamental strings but branes wrapped on a 4-cycle of the Calabi-Yau manifold. Towards the end of inflation the volume of the 4-cycle becomes very small as the Calabi-Yau volume shrinks to a very small size, endowing the corresponding strings with a low tension. There is an extensive production of this type of strings (a similar situation for nonperturbative strings obtained by wrapping D3-branes on shrinking 2-cycles has been discussed in [71] and references therein). A very rough estimate shows that the effective tension of a string obtained by wrapping a brane on a nontrivial cycle has to satisfy

$$\sqrt{\mu_{\text{string}}} \leq \frac{1}{20} M_{\text{Pl}}, \tag{6.3}$$

in order to lead to a massive string production. This bound can be easily satisfied for the case of an $M5_{\|}$-brane. In this case the effective string tension is given by

$$\mu_{M5_{\|}} = \tau_{M5}(1 - x^{11}_{M5}/L_c)^{2/3} V_{\Sigma_4}. \tag{6.4}$$

This expression makes it clear that we can easily satisfy the bound (6.3) by being close enough to the hidden boundary where the warp factor can be made arbitrarily small. As a result, tensionless cosmic strings will be produced on the hidden boundary. Even though the tensionless strings are produced on the hidden boundary, they still would have an effect on our visible universe since they interact gravitationally. These strings would then represent an interesting new dark matter candidate (for their detection via gravitational lensing see e.g. [75–77]) next to other potential dark matter residing on the hidden boundary [78]. One final remark on the stability of the pair produced strings is in order. For the pair produced strings to be observed, it is important that they stay around long enough and do not annihilate shortly after being pair produced. Even though annihilation and decay of strings are still poorly understood, one can make several arguments in favor of the stability and observability of the strings being pair produced. One qualitative argument is based on the dimensionality of the string world sheet and was used many years ago by Brandenberger and Vafa to argue that our world is 4 dimensional [79]. We could argue that the odds for an infinitely extended string pair to meet once produced are pretty small, as the world sheet of a string is 2 dimensional and two strings would only collide at an instant. A generalization of this idea, worked out more recently in a paper by Randall and Karch [80], would allow us to exclude the production of higher dimensional branes, as the odds for such branes to meet and annihilate are much higher. It would be interesting to work out the details of this higher order annihilation process more precisely. We hope to report on this and on the production rate calculation of the cosmic string candidate presented in this paper elsewhere.

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