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SPACE-TIME SUPERSYMMETRY BREAKING IN HETEROTIC STRING THEORY

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A B S T R A C T

A simple mechanism for space-time supersymmetry breaking in the ten-dimensional heterotic string theory is proposed. We present a heterotic string model with a hidden two-dimensional sigma-model sector which can induce desirably small supersymmetry breaking, without upsetting the zero value of the cosmological constant, through a topological instanton effect due to an Abelian gauge field on the string world-sheet. We find that the consistency condition of the gravitino field equation is satisfied for some configurations even after supersymmetry breaking.

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1. INTRODUCTION. There have been several long-standing tasks for the ten-dimensional (heterotic) superstring theories [1] to be acceptable as genuine theories unifying all the interactions. One of them is to show the all-order finiteness in a mathematically rigorous way. Another one is to show how space-time supersymmetry is broken keeping the huge mass hierarchy between the Planck mass scale (10^{19} GeV) and the W-boson mass scale (10^2 GeV). Since gravitational interactions are essential in string theories, we also have to keep the zero (or extremely small) value of the cosmological constant to all orders even after supersymmetry breaking. Unless we solve this problem, superstring theories can never be realistic theories describing low energy physics, even though grand unification is successful [2].

In this paper we try to take a step towards solving the second problem, i.e., breaking space-time supersymmetry at low energy, while maintaining the zero cosmological constant. We present a simple toy-model, which contains the ten-dimensional heterotic string theory [1], not only with the good phenomenological features of the heterotic strings before the supersymmetry breaking, but also with a reasonably small amount of supersymmetry breaking. Even though we do not claim that it is the unique ultimate theory, we hope that our model can provide a new method for supersymmetry breaking at low energy, in any superstring theory even in lower dimensions.

In recent papers by Dine et al. [3], it is claimed that world-sheet holomorphic instantons associated with the Wess-Zumino term can break four-dimensional ($D = 4$) supersymmetry. Unfortunately, in that approach the re-

sultant potential in the case of the (2,0) heterotic sigma-model cannot have a stable vacuum with the zero cosmological constant. Another interesting approach is to introduce a new type of symmetry called Atkin-Lehner symmetry [4], which seems attractive as an alternative of supersymmetry. However, there is a question about higher-loop effects, and the protection of the gauge hierarchy with the desirably small (but non-zero) mass scale is not quite obvious. In the case of introducing the Fayet-Iliopoulos D-term [5], an extra U(1) field is needed, and moreover the supersymmetry breaking scale becomes undesirably as heavy as the Planck mass, and a huge cosmological constant may be induced.

Our model instead makes use of two-dimensional vector supermultiplets, which provide an alternative topological instanton effect leading to space-time supersymmetry breaking with stable vacuum and zero cosmological constant.

2. FIELD CONTENT AND LAGRANGIAN. We first fix the field content of our theory, restricted by various anomaly cancellation requirements in two dimensions. We utilize some vector supermultiplets coupled to (1,0) heterotic sigma-models in $D = 2$ [6] in order to induce instanton effects on the string world sheet.

Suppose we have V copies of U(1) supermultiplets $(A_i^{(r)}, \lambda_L^{(r)})$, $((r) = 1, \dots, V)$, S extra real scalar multiplets $(\phi^{\mu'}, \chi_R^{\mu'})$, $(\mu' = 1, \dots, S)$, U extra unidexterous fermions $\psi_L^{A'}$, $(A' = 1, \dots, U)$, in addition to the usual $E_8 \times E_8$, $D = 10$ heterotic sigma-model fields: (ϕ^μ, χ_R^μ) , $(\mu = 1, \dots, 10)$, and ψ_L^A , $(A = 1, \dots, 32)$. Conformal invariance forbids kinetic terms for the vector multiplets. Although they do not contribute directly to the anomalies, their ghosts do. In $D = 2$ one expects to gauge-

fix to zero both components of a vector field, due to the presence of two kinds of $U(1)$ symmetries: $U(1)_V \times U(1)_A$ [6,7,8], and we have maximally two pairs of Faddeev-Popov ghosts. Performing a supergraph calculation for the contributions from the vector multiplets to the conformal and gravitational anomalies as in Ref.[7], we find that these anomalies are cancelled, when

$$\left[10 + \frac{32}{2} - 26 + S + \frac{U}{2} - \frac{a}{2} V\right] (D_+ H_+^{\bar{+}}) \square^{-1} \partial_-^4 H_+^{\bar{+}} = 0, \quad (2.1)$$

$$(10 - 10 - S) (D_+ H_-^{\bar{+}}) \square^{-1} \partial_+^3 H_-^{\bar{+}} = 0, \quad (2.2)$$

where we use the notation of Ref.[7]. Obviously the solution for cancellation of anomalies is

$$S = 0, \quad U = a V. \quad (2.3)$$

The constant a can be 2, 1 or 0 denoting the number of the gauge-fixed $U(1)$ groups out of $U(1)_V \times U(1)_A$. For example, if we gauge-fix only $U(1)_V$, we have $a = 1$. For simplicity we choose

$$U = a, \quad V = 1, \quad S = 0. \quad (2.4)$$

We mention that our choice is just one example among anomaly-free cases. If we allow minimal couplings to the extra unidexterous fermion multiplet, we generally encounter an $U(1)_V$ (and/or $U(1)_A$) anomaly [9]. This is because the minimal couplings to the right-handed fermions χ_R^μ , which could cancel this $U(1)$ anomaly in turn cause some undesirable mixing between the extra multiplets and the usual heterotic string part, and may upset the desirable structures of the $D = 10$ heterotic string. Therefore we exclude these minimal couplings to avoid $U(1)$ anomalies.

From these considerations, we are led to the Lagrangian,

$$\begin{aligned}
 e^{-1} L_{D=2} = & -\frac{1}{2} (g^{ij} g_{\mu\nu} + \alpha \varepsilon^{ij} B_{\mu\nu}) \partial_i \phi^\mu \partial_j \phi^\nu \\
 & - \frac{i}{2} g_{\mu\nu} \bar{\chi}^\mu \gamma^i [\partial_i \chi^\nu + \Gamma_{\rho\sigma}^\nu (\partial_i \phi^\rho) \chi^\sigma] - i g_{\mu\nu} \bar{\psi}_i \gamma^j \gamma^i \chi^\mu \partial_j \phi^\nu \\
 & - \frac{1}{6} \alpha T_{\mu\nu\rho} \bar{\psi}_i \chi^\mu \bar{\chi}^\nu \gamma^i \chi^\rho - \frac{i}{2} G_{AB}(\phi) \bar{\psi}^A \gamma^i [\partial_i \psi^B + A_\mu{}^B{}_C (\partial_i \phi^\mu) \psi^C] \\
 & - \frac{1}{8} \hat{F}_{\mu\nu AB} \bar{\chi}^\mu \gamma^i \chi^\nu \bar{\psi}^A \gamma_i \psi^B \\
 & - \frac{i}{2} \sum_{A'=1}^a \bar{\psi}^{A'} \gamma^i \partial_i \psi^{A'} + g' e^{-1} \varepsilon^{ij} F_{ij} \Theta(\phi^\mu) - 4i g' \bar{\lambda} \chi^\mu \partial_\mu \Theta(\phi^\nu). \quad (2.5)
 \end{aligned}$$

As mentioned above, the extra vector multiplet has no kinetic term. Such a term would spoil the conformal invariance of the classical lagrangian. (The $\varepsilon^{ij} F_{ij}$ -term was not given in Ref.[6], but its superspace form is provided in Refs.[7,8].) At this stage $\Theta(\phi^\mu)$ can be any scalar function of ϕ^μ in general. However, the field equation of the vector field A_i would require that the Θ is a world-sheet constant. Therefore it can only depend on ϕ_B^μ , the zero mode of ϕ^μ . Eventually the εF -term is a total divergence. The Θ 's is a world-sheet constant, and there is no contribution from the vector field, except instanton-type solutions. Therefore we should choose the case $a = 0$, because the εF -term is a total divergence and we do not need any gauge-fixing with Faddeev-Popov ghosts. For the two-dimensional constant $\Theta = \Theta(\phi_B^\mu) = \Theta_B$, the penultimate term in eq.(2.5) can be interpreted as accounting for the existence of various Θ -vacua. Since the vector field A_i has no degrees of freedom, the path-integration over it reduces to the sum over various Θ -vacua.

To see the absence of tachyons in our system, we have to study the mass spectra, which require field expansions satisfying the field equations

[1]. As usual, we set all the background fields in the target space to zero. Let us consider the field equations for the extra multiplet of $a = 0$ case. (We call this the hidden sector.) Let L_H be a lagrangian for the hidden sector: $e^{-1}L_H = g'e^{-1}\varepsilon^{ij}F_{ij}\theta_B$. Notice the fact that the εF -term does not contribute to the energy-momentum tensor, and our hidden sector decouples from the rest. This means that when we construct the Virasoro operators, we do not have to consider the effect of the (A_i, λ_L) -multiplet.

It is now easy to see that the mass spectrum of our system is exactly the same as in the heterotic superstring theory in $D = 10$, and therefore, we have no tachyons [1]. The one-loop modular invariance is also guaranteed, as in the heterotic string theory [1]; The partition function A_F for the fermions can be written as $A_F = A_R A_L$,

$$A_R = \sum_{\alpha} A_R^{\alpha}(\alpha) = \left[\theta_2^4(0|\tau) - \theta_3^4(0|\tau) + \theta_4^4(0|\tau) \right] / \eta^4(\tau) \equiv 0, \quad (2.6)$$

where A_R (or A_L) is for the right (or left) moving fermions. The α is for the boundary conditions on the torus. Of course the M -point ($M \geq 4$) 1-loop string amplitudes are exactly the same as the heterotic string theory; our system is also modular invariant. If we follow the ordinary argument for higher-genus cases [10], we see that the modular invariance at higher-loops is also guaranteed in our model.

3. INSTANTON EFFECTS AND SUPERSYMMETRY BREAKING. So far we have seen that our extra vector multiplet has no effect on the $D = 10$ in a flat background. If we include the $D = 10$ background and consider the $D = 2$ instanton effects as well, this situation drastically changes.

If we allow the nontrivial instanton solution on the world sheet:

$$\int d^2x \varepsilon^{ij} F_{ij}^N = 2\pi N, \quad (N = 1, 2, \dots), \quad (3.1)$$

then our εF -term gives an N -dependent contribution to the $D = 10$ effective action [11], just as the $R^{(2)}\Phi$ -term provides the contribution related to the Euler number of the world sheet. We have the $D = 10$ effective action

$$\begin{aligned} \Gamma_{D=10} = \sum_{\text{topologies}} \int d^{10}\phi_B \int [dg_{ij}] [d\xi^i] \dots \exp \left[-\frac{1}{2\pi} \int d^2x e \right. \\ \left. \times \left\{ -\frac{1}{2} g^{ij} g_{\mu\nu} \partial_i \phi^\mu \partial_j \phi^\nu + \dots + \frac{1}{2} R^{(2)} \Phi + g' \varepsilon^{ij} F_{ij} \Theta(\phi_B^\mu) \right\} \right], \end{aligned} \quad (3.2)$$

where $\phi^\mu = \phi_B^\mu + \xi^\mu$, and $\phi_B^\mu = \text{const.}$ is the "center of mass" collective coordinate [11]. The εF -term contributes $\exp(-Ng'\Theta)$ to $\Gamma_{D=10}$ as a universal factor, as in the case of $\exp[2(h-1)\Phi(\phi_B^\mu)]$ from the $R^{(2)}\Phi$ -term. (Here h is a genus.) Notice that even though our εF -term is just a total divergence in $D = 2$, it has a physical significance in $D = 10$, as the $R^{(2)}\Phi(\phi_B)$ -term does. The term $\varepsilon^{ij} F_{ij} \Theta$ is the only N -dependent term in the exponent in eq.(3.2) for the constant Θ . (The A_i field is not dynamical, and we do not have to consider its quantum fluctuations.) Therefore,

$$\Gamma_{D=10} = \sum_{h=0}^{\infty} \sum_{N=0}^{\infty} \int d^{10}\phi_B e^{2(h-1)\Phi_B - Ng'\Theta_B} \tilde{L}[\partial\Phi_B, g_{\mu\nu}, \dots] \quad (3.3)$$

$$\equiv \int d^{10}\phi_B \sum_{h=0}^{\infty} \sum_{N=0}^{\infty} L_h^N \equiv \int d^{10}\phi_B \sum_{h=0}^N L_h, \quad (3.4)$$

where dots denote other background fields in $D = 10$, and $\Phi_B \equiv \Phi(\phi_B^\mu)$, $\partial\Phi_B \equiv \partial\Phi_B/\partial\phi_B$. The Φ 's appears in $\tilde{L}$ only through its derivatives, and $\Gamma_{D=10}$ is N -independent.

We now restrict ourselves to the tree-level ($h = 0$). In general, θ is a constant on a given world-sheet, therefore it can be any scalar function of ϕ_B^μ . Thus θ_B 's can be also functional of $D = 10$ massless background fields, such as the dilaton. We have now

$$L_{h=0} \equiv F(\theta_B) e^{-2\Phi_B} \tilde{L}[\partial\Phi_B, g_{\mu\nu}, \dots], \quad (3.5)$$

where $F(x) = 1/(1 - e^{-g'x})$.

The effect of our g' -term is to spoil the global Weyl-weight [12]: $2 - D = -8$ of $L_{h=0}$. This in turn signals breaking of supersymmetry, which can be seen as follows: Suppose our system had Poincaré supersymmetry, in which case it could be always embedded into off-shell conformal supergravity with $128 + 128$ degrees of freedom [12]. Even after imposing constraints to recover Poincaré supergravity, there would still remain a global Weyl covariance. This would be manifest especially in the Yang-Mills kinetic term in $L_{h=0}$. (Recall that the Yang-Mills field always has zero Weyl weight [12].) However, in our case the global Weyl covariance is impossible because of the anomalous dependence on the compensator field (dilaton) through the factor $F(\theta_B)$, which generally contains the dilaton. Therefore with a dilaton-dependent functional $F(\theta_B)$, our $D = 10$ massless field system cannot be regarded as coming from any off-shell conformal supergravity in $D = 10$, and this implies the breaking of local supersymmetry. (We are of course relying on the empirical fact that any Poincaré supergravity can be embedded in off-shell conformal supergravity, and on the result that the global Weyl-covariance remains after conformal breaking constraints are imposed.)

We view our model as providing a mechanism for explicit breaking of $D = 10$ supersymmetry. However, since the breaking is induced at the world-sheet level, and modular invariance and other string features are maintained, it seems reasonable to assume that it will not destroy good string properties such as finiteness, so that the supersymmetry breaking is "soft".

Essentially the same argument holds for higher-genus ($h \geq 2$). This is because the N -dependence is only in the F_{ij} -linear term in the two-dimensional lagrangian (2.5), and the $D = 10$ supersymmetric effective action $L_h^{N=0}$ has the zero cosmological constant by the non-renormalization theorem [3,13]; eventually, the N -summation gives only the supersymmetry-breaking effect by supplying an overall factor in front of the potential.

4. CONSISTENCY OF GRAVITINO FIELD EQUATION. We next consider what kind of functional Θ can be actually introduced in a consistent way. In models with explicit breaking of local supersymmetry, there generally arise some inconsistencies. This is because there are strong restrictions by the consistency of the gravitino field equation [14]. To find the most general resolution to this problem is quite difficult, but as is shown below, there is an explicit example of $\Theta(\phi_B^\mu)$ to satisfy this consistency.

Let L_0 be the string-corrected $D = 10$, $N = 1$ supergravity lagrangian at $h = 0$ and $g'N = 0$, and let L be our total lagrangian at $h = 0$:

$$L_0 \equiv L_{h=0}^{N=0} = e^{-2\Phi_B} \tilde{L}[\partial\Phi_B, g_{\mu\nu}, \dots], \quad (4.1)$$

$$L \equiv L_{h=0} = F(\Theta_B) L_0. \quad (4.2)$$

Let us also assume that F is a function only of $e^{-1}L_0$: $F = G(e^{-1}L_0)$. Obviously $\partial L_0 / \partial \Phi = -2L_0$. However, it is easy to show that by adding an appropriate total divergence to L_0 , we can have $\delta L_0^{\text{new}} / \delta \Phi = -2L_0^{\text{new}}$ for the new L_0 . Therefore for this L_0^{new} , the dilaton field equation is completely equivalent to $L_0^{\text{new}} = 0$. Henceforth, we always use L_0^{new} instead of the original L_0 by omitting the superscript "new" for convenience's sake.

The consistency condition [14] for our gravitino field equation $R^\mu \equiv \delta L / \delta \psi_\mu = 0$ is that the equation $D_\mu R^\mu = 0$ should be satisfied at least for some non-trivial (dynamical) solutions of our field equations. Our field equation for an arbitrary field φ^I is¹

$$\begin{aligned} \frac{\delta L}{\delta \varphi^I} &= \frac{\partial G}{\partial \varphi^I} L_0 + G \frac{\partial L_0}{\partial \varphi^I} - \partial_\mu \left[\frac{\partial G}{\partial (\partial_\mu \varphi^I)} L_0 + G \frac{\partial L_0}{\partial (\partial_\mu \varphi^I)} \right] \\ &\quad + \partial_\mu \partial_\nu \left[\frac{\partial G}{\partial (\partial_\mu \partial_\nu \varphi^I)} L_0 + G \frac{\partial L_0}{\partial (\partial_\mu \partial_\nu \varphi^I)} \right] - \dots \\ &= G \frac{\delta L_0}{\delta \varphi^I} + (\text{terms with } L_0 \text{ or } \partial_{\mu_1} \dots \partial_{\mu_n} L_0). \end{aligned} \quad (4.3)$$

Since G is only a function of $e^{-1}L_0$, all the terms besides $G \delta L_0 / \delta \varphi^I$ in eq.(4.3) always contain L_0 itself or its derivatives $\partial_{\mu_1} \dots \partial_{\mu_n} L_0$. This structure ensures that there are solutions satisfying the unbroken supersymmetric field equations $\delta L_0 / \delta \varphi^I = 0$, $L_0 = 0$. This aspect is highly non-trivial, because, solutions of $\delta L_0 / \delta \varphi^I = 0$ are still elements of the whole set of solutions of the field equations with broken supersymmetry, which are much more complicated and general. Moreover, it is now easy to see the

¹ Since L_0 itself in general contains higher-order derivatives [1], we need to consider higher-derivative-terms in the Euler derivative.

satisfaction of our consistency condition:

$$\begin{aligned}
 D_\mu R^\mu &= D_\mu \left[G \frac{\delta L_0}{\delta \psi_\mu} + (L_0 \text{ or } \partial_{\mu_1} \dots \partial_{\mu_n} L_0) \right] \\
 &= (\partial_\mu G) \frac{\delta L_0}{\delta \psi_\mu} + G D_\mu \left(\frac{\delta L_0}{\delta \psi_\mu} \right) + (L_0 \text{ or } \partial_{\mu_1} \dots \partial_{\mu_n} L_0) \\
 &= 0.
 \end{aligned} \tag{4.4}$$

Since we already know that $D_\mu(\delta L_0/\delta \psi_\mu) = 0$ can be satisfied by the solutions of $\delta L_0/\delta \varphi^I = 0$, $L_0 = 0$ (exactly supersymmetric case), it is obvious that eq.(4.4) is really satisfied by the solutions of $\delta L_0/\delta \varphi^I = 0$, and thus the consistency condition is satisfied at least for non-trivial solutions.

Remarkably enough, the cosmological constant of our system remains exactly zero for any solution, due to the fact that

$$\frac{\partial L}{\partial \Phi} = \left(\frac{\partial G}{\partial \Phi} - 2 G \right) G^{-1} L \propto L. \tag{4.5}$$

(Here we need only the partial derivative instead of Euler derivatives to see the minimum of the potential.) In other words, our cosmological constant is always self-fine-tuned to be zero!

The degeneracy of the usual flat supersymmetric vacua is now enlarged to include non-supersymmetric vacua. Since the breaking of the supersymmetry is now a matter of vacuum configurations, we can regard our breaking mechanism as a generalized super-Higgs effect. In other words, our supersymmetry breaking is also "spontaneous".

Recall that our parameter g' controls the magnitude of our supersymmetry breaking. This is not like the situation in Ref.[3], because our g' is not quantized because we directly work in 10 dimensional spacetime rather than in $K4 \times K6$. Since this breaking scale is an exponential func-

tion, it can be made desirably small (e.g., $e^{-g'\theta} \approx 10^{-6}$ or 10^{-16} .) by choosing the $O(1)$ value of g' . It is to be stressed that both the cosmological constant and the small supersymmetry breaking are satisfied, keeping the consistency of the gravitino field equation.

5. DISCUSSIONS AND REMARKS. In this paper we have presented a simple extension of the heterotic string theory keeping the good properties of the heterotic string such as modular invariance and the absence of tachyons. Our model can have a small amount of the $D = 10$, $N = 1$ supersymmetry breaking, with the zero cosmological constant, satisfying also the consistency condition of the gravitino field equation. Since the sigma-model structure of $(1,0)$ heterotic strings in general admits an arbitrary scalar functional θ in the $\varepsilon^{ij}F_{ij}$ -term, we expect that similar mechanisms of space-time supersymmetry-breaking can work in any other heterotic string theories in lower dimensions [15] based on $(1,0)$ heterotic sigma-models.

Our example of the field equation system from $L = G(e^{-1}L_0)L_0$ has quite an interesting structure. The system contains the solutions with unbroken supersymmetry as well as solutions with broken supersymmetry, maintaining also the consistency condition and the zero cosmological constant. This implies that our vacuum structure has a huge degeneracy which is bigger than that of the usual flat supersymmetric vacua. Considering the fact that L_0 itself contains higher-order derivatives, it is not unnatural to think of such a higher-derivative lagrangian as $L = G(e^{-1}L_0)L_0$, which accordingly results in the huge degeneracy of the vacuum. It is no wonder to encounter this degeneracy under higher-derivatives and non-perturbative effects of superstrings.

Our $\varepsilon^{ij}F_{ij}$ -type term, which is a topological analogue of the $R^{(2)}\Phi$ -term, plays an important role for the supersymmetry breaking with the vanishing cosmological constant. The cosmological constant after the supersymmetry-breaking at higher-loops also vanishes, provided that the exactly supersymmetric effective potential has the zero cosmological constant guaranteed by the non-renormalization theorem [13]. Now the long-standing question "Why is the cosmological constant always self-fine-tuned to be zero at tree-level as well as at higher-orders even after the supersymmetry breaking?" is well answered by our mechanism. As for cosmological dynamical behaviours of our system, we suspect that our functional Θ appearing in the exponent with the dilaton plays a very special role very similar to that of the usual dilaton in the supergravity cosmological models studied in Ref.[16]. The Friedmann universe will be automatically realized at the final stage of the universe [16]. We conclude by mentioning that our model may well be even a peculiar example with the Atkin-Lehner symmetry [4].

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