The Non-Perturbative $SO(32)$ Heterotic String

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ABSTRACT

The $SO(32)$ heterotic string can be obtained from the type IIB string by gauging a discrete symmetry that acts as $(-1)^{F_L}$ on the perturbative string states and reverses the parity of the D-string. Consistency requires the presence of 32 NS 9-branes – the S-duals of D9-branes – which give $SO(32)$ Chan-Paton factors to open D-strings. At finite string coupling, there are $SO(32)$ charges tethered to the heterotic string world-sheet by open D-strings. At zero-coupling, the D-string tension becomes infinite and the $SO(32)$ charges are pulled onto the world-sheet, and give the usual $SO(32)$ world-sheet currents of the heterotic string.
The perturbative type I superstring can be obtained from the perturbative type IIB superstring by orientifolding by the action of the world-sheet parity operator $\Omega$ [1,2]. However, it seems that $\Omega$ extends to a symmetry of the full non-perturbative IIB theory; indeed, it is straightforward to extend its action to BPS states and to the dual IIB string theory that emerges in the strong coupling limit, where it acts through a perturbative symmetry $\tilde{\Omega}$ of the dual theory. If $\Omega$ does extend to such a non-perturbative symmetry, we can consider modding out the IIB theory by this symmetry for any value of the string coupling. In particular, in the strong coupling limit, this corresponds to modding out the dual IIB string theory by $\tilde{\Omega}$, and this should give the $SO(32)$ heterotic string, as this is the strong coupling limit of the type I string [3-6]. This led to the conjecture of [7], that the $SO(32)$ heterotic string can be obtained from the type IIB superstring by modding out by the action of the symmetry $\tilde{\Omega}$. The purpose of this paper is to investigate this construction further, and in particular show how the gauge sector of the heterotic string emerges.

In the orientifold construction of the type I superstring, it is essential to add 32 D9-branes which provide the Chan-Paton factors for the open strings. Extrapolating to strong coupling, the D9-branes are replaced by their S-duals, which are the NS-NS 9-branes proposed in [7]. Then 32 NS-NS 9-branes are required in the construction of the $SO(32)$ heterotic string from the IIB string [7] and should play a key role in the construction of the gauge sector. We shall show that the 9-branes indeed give rise to the gauge sector. Moreover, this extends to a construction of the non-perturbative heterotic string and in particular we will find that at finite string coupling, the $SO(32)$ charges are no longer confined to the string world-sheet but are tethered to it by strings, and as these strings can break, the charges can escape.

The type IIB string theory has an $SL(2,\mathbb{Z})$ U-duality symmetry [8], and the massless bosonic fields are $g_{MN}, B^{1}_{MN}, \Phi$ in the NS-NS sector and $D_{MNPQ}, B^{2}_{MN}, \chi$ in the RR sector. The 2-form fields $B^{1}_{MN}$ transform as doublets under $SL(2)$, while $SL(2)$ acts on $\lambda = \chi + ie^{-\Phi}$ through fractional linear transformations. There is also a non-dynamical RR 10-form potential $A^{2}_{M...N}$ that couples to D9-branes [6].
must fit into an $SL(2)$ doublet also [7], and its partner is a NS-NS 10-form potential $A_{1M...N}^i$. The theory has $Dp$-branes for $p = 1, 3, 5, 7, 9$ [9] and it is interesting to ask how these transform under $SL(2, \mathbb{Z})$. The $D3$-brane is invariant, but acting on the D1, D5 and D9 branes generates $(p, q)$-branes for all co-prime integers $(p, q)$, as can be seen from the superalgebra [7]. Of these 1-branes, 5-branes and 9-branes, the $(0,1)$-branes are D-branes, the $(1,0)$-branes couple to fields in the NS-NS sector, while the $(p, q)$ branes are bound states of $p (1,0)$ branes and $q (0,1)$ branes (for co-prime integers $(p, q)$). The 1-branes and 5-branes couple to the 2-form fields $B_{MN}^i$ and the 9-branes can be thought of as coupling to 10-form potentials $A_{M...N}^i$.

The 7-brane is more subtle, as it couples to the scalar fields which transform non-linearly under $SL(2, \mathbb{Z})$. The solutions of [10], in which the axion ansatz involves the modular invariant $j$-function, are $SL(2, \mathbb{Z})$ invariant and lead to singlet 7-brane charges. Acting with $SL(2, \mathbb{Z})$ on a 7-brane leaves its charge $Z_{i_1...i_7}$ invariant, but changes the $SL(2, \mathbb{Z})$ monodromy and the couplings to strings and 5-branes. These branes are characterised by two integers [11] and are obtained from the $(0,1)$ D7-brane of [10] by an $SL(2, \mathbb{Z})$ transformation; we will refer here to the 7-branes on which a $(0,1)$ string can end as $(1,0)$ 7-branes.

At weak coupling, $g \equiv <e^{-\Phi}> \approx 0$, the perturbative states are described by the NS-NS or $(1,0)$ string while all the other branes are non-perturbative [12]. The perturbative theory is formulated as the usual type IIB superstring theory, with a topological expansion in terms of the genus of the world-sheet of the NS-NS string. At strong coupling, however, it is the RR or $(0,1)$ string that gives the states that are perturbative in an expansion in $\tilde{g} = 1/g$. The self-duality of the theory implies that the strong-coupling theory is again a type IIB string theory, but now the formulation should be in terms of the world-sheet of the $(0,1)$ string. The perturbation theory in $\tilde{g}$ is a sum over $(0,1)$ string world-sheets with the power of $\tilde{g}$ corresponding to the genus of the world-sheet. The dual perturbative string theory has left and right movers which each decompose into an NS and an R sector. For example, the fundamental string of the weakly coupled ($g$-perturbative) theory couples to the 2-form in the NS-NS sector but becomes the D-string of the dual
theory, coupling to the 2-form in the RR sector of the $\tilde{g}$-perturbative theory.

At weak coupling, the (1,0) string can end on the D-branes carrying RR charge, which are the 3-brane, and the (0,1) $p$-branes with $p = 1, 5, 7, 9$. In the strongly coupled theory formulated in terms of the world-sheet of the (0,1) string, the (0,1) string, the (0,1) 5-brane and the (0,1) 9-brane carry charge that appears in the NS-NS sector of the dual (0,1) string while the new D-branes with density proportional to $1/\tilde{g}$ on which the (0,1) string can end are the 3-brane, and the (1,0) $p$-branes with $p = 1, 5, 7, 9$. These dual D-branes all carry charge which occurs in the RR sector of the (0,1) string. This structure is obtained by acting with the $SL(2,\mathbb{Z})$ transformation

$$S = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$$ (1)

which interchanges strong and weak coupling. For example, this takes a (1,0) string ending on (0,1) strings or 5-branes to a (0,1) string ending on (1,0) strings or 5-branes.

The perturbative type IIB theory is formulated in terms of a fundamental (1,0) string and has a symmetry $\Omega$ which reverses the parity of the (1,0) world-sheet. The dilaton is invariant under $\Omega$, which is a perturbative symmetry (i.e. it takes perturbative states to perturbative states). Of the massless bosonic fields, the ones that are invariant under $\Omega$ are $g_{MN}, B^2_{MN}, \Phi, A^2_{M...N}$, while the others are odd, and this tells us how it acts on BPS states: the (0,1) string, 5-brane and 9-brane are invariant while for the (1,0) string, the 2-form $B^1_{MN}$ to which it couples is odd, but the Wess-Zumino term in the world-sheet action is invariant as the world-sheet volume form is odd under world-sheet parity $\Omega$. Similarly, the (1,0) 5-brane and 9-brane together with the 7-brane and 3-brane couple to $p + 1$ form gauge potentials that are odd under $\Omega$, and so their world-volume volume-forms must also be odd under $\Omega$, so that $\Omega$ must act as an orientation-reversing world-volume parity operator on these branes.

The branes that couple only to the invariant fields are invariant, so that the
only \((p, q)\) D-branes that are invariant are the D1,D5,D9 branes with charges \((0,1)\). This tells us how to extend the action of \(\Omega\) to the BPS brane sector, and it will be assumed that the perturbative symmetry extends to a symmetry of the full non-perturbative type IIB string (which will also be denoted \(\Omega\)). This is the key assumption; if \(\Omega\) did not extend to a symmetry of the full non-perturbative type IIB theory, then orientifolding by \(\Omega\) would be problematic even for small but finite string coupling, whereas it is believed that the type I string extends to a consistent non-perturbative theory, and this assumption is implicitly made in many discussions of string dualities. The existence of such a non-perturbative symmetry can be ‘derived’ from M-theory: the non-perturbative IIB theory arises from M-theory compactified on a 2-torus in the limit in which both radii tend to zero [13] and the symmetry that provides the non-perturbative version of \(\Omega\) is given by the IIB theory limit of the \(\mathbb{Z}_2\) M-theory symmetry used by Hořava and Witten [14].

Then

\[
\tilde{\Omega} = S\Omega S^{-1}
\]

is also a symmetry of the full IIB theory, where \(S\) is the \(SL(2,\mathbb{Z})\) transformation (1) interchanging weak and strong coupling. However, \(\tilde{\Omega}\) leaves the dilaton invariant, and is a symmetry order by order in IIB string perturbation theory. The weakly coupled type IIB string has a perturbative symmetry \((-1)^{F_L}\), where \(F_L\) is the left-handed fermion number of the conventional world-sheet formulation, and as \(\tilde{\Omega}\) and \((-1)^{F_L}\) act in exactly the same way on perturbative states, \(\tilde{\Omega}\) can be thought of as a non-perturbative extension of \((-1)^{F_L}\) [7]. The massless bosonic fields that are invariant under \(\tilde{\Omega}\) are \(g_{MN}, B^1_{MN}, \Phi, A^1_{M...N}\), which are precisely the NS-NS fields, so that D-branes are odd and NS or \((1,0)\) strings, 5-branes and 9-branes are even under \(\tilde{\Omega}\). In the same way that \(\Omega\) inverts the parity of the \((1,0)\) string world-sheet, together with the world-volume orientations of the \((1,0)\) 5-brane and 9-brane and the 7-brane and 3-brane, \(\tilde{\Omega}\) inverts the parity of the \((0,1)\) string world-sheet, and also inverts the world-volume orientations of the \((0,1)\) 5-brane and 9-brane and the 7-brane and 3-brane.
The orientifold of the weakly coupled IIB theory constructed using $\Omega$ gives the weakly coupled type I theory [1,2]. The orientifolding can be thought of as introducing an orientifold 9-plane, and 32 D9-branes must be added to cancel the anomalies and divergencies introduced by the orientifold 9-plane. The invariant sector is formulated in terms of unoriented closed type I strings, with massless bosonic fields $g_{MN}, B^2_{MN}, \Phi$. In addition there is an open string sector, which can (in some ways) be thought of as a twisted sector, with $SO(32)$ Chan-Paton factors arising from 32 RR 9-branes; the Chan-Paton factor labels which 9-brane the string ends on [6]. The RR string, 5-brane and 9-brane are invariant under $\Omega$, and so survive the projection. The closed fundamental IIB strings are BPS, and as $\Omega$ acts as world-sheet parity on these, they become the nonoriented fundamental type I strings, which are no longer BPS and can break. As well as non-oriented closed (1,0) strings, there should also be non-oriented (1,0) 5-branes, 3-branes and 7-branes (and, in principle, (1,0) 9-branes) in the theory, as $\Omega$ acts as an orientation-reversing parity operator on their world-volumes.

An isomorphic construction emerges if we instead mod out the strongly coupled IIB string using $\tilde{\Omega}$. Indeed, the strongly coupled IIB theory is formulated as a world-sheet theory of the (0,1) string, which is a perturbation theory in $\tilde{g} = 1/g$, and $\tilde{\Omega}$ reverses the parity of the (0,1) string world-sheet. The resulting type I theory is one in which the fundamental open and closed strings are (0,1) strings which can end on the (1,0) strings, 5-branes and 9-branes. Consistency requires 32 (1,0) 9-branes (which can be thought of as cancelling out the effects of a (1,0) orientifold 9-plane), and these give $SO(32)$ Chan-Paton factors to the open (0,1) strings. This is the standard type I superstring, but embedded differently in the type IIB theory. If this extrapolates to all values of the coupling, then continuing back to weak coupling, we find that modding out weakly-coupled type IIB string by $\tilde{\Omega}$ in the presence of 32 (1,0) 9-branes should give the weakly coupled $SO(32)$ heterotic string. This construction of the $SO(32)$ heterotic string from the IIB string modded out by $\tilde{\Omega}$ with 32 9-branes can be understood as a particular limit of the Hořava-Witten construction. In [14], M-theory on $T^2$ is modded out by a
\( \mathbb{Z}_2 \) symmetry to give M-theory on a cylinder \( S^1 \times S^1 / \mathbb{Z}_2 \). In the limit in which the radius \( R \) of the \( S^1 \) and the length \( L \) of the \( S^1 / \mathbb{Z}_2 \) both tend to zero, with the limiting form of \( L/R \) small, the M-theory on \( T^2 \) becomes the IIB string with coupling \( L/R \), the \( \mathbb{Z}_2 \) symmetry becomes \( \tilde{\Omega} \) and the resulting theory is the \( SO(32) \) heterotic string. Similarly, if \( L/R \) is large, this reduces to the orientifold construction of the type I string; see [15] for details.

This construction implies that the gauge structure of the heterotic string must arise from the 32 (1,0) 9-branes. We will now show how this comes about, and verify that the perturbative heterotic string indeed emerges from this construction.

Consider first the D-string in the weakly-coupled type I string theory. The zero-mode structure of the supergravity solution corresponding to the D-string was studied in [4,5] and shown to give precisely the world-sheet structure of the heterotic string. The excitation spectrum of the D-string can be calculated by quantizing the open string allowing Dirichlet (D) boundary conditions with the fundamental strings ending on the D-string, as well as Neumann (N) boundary conditions [6]. There are three sectors: (i) the NN sector in which both ends of the string satisfy Neumann boundary conditions (or, equivalently, end on the D9-branes) and carry \( SO(32) \) Chan-Paton factors, giving the usual open strings of type I string theory; (ii) the DD sector in which both ends of the string lie on the D-string has a massless sector which is a world-sheet theory on the D-string, consisting of 8 scalars and 8 right-handed Majorana-Weyl world-sheet fermions; (iii) the DN sector in which one end of the string lies on the D-string and the other satisfies Neumann boundary conditions and carries an \( SO(32) \) Chan-Paton factor has a massless sector which is again a theory on the D-string world-sheet, this time consisting of 32 left-handed Majorana-Weyl world-sheet fermions transforming as a \( 32 \) of \( SO(32) \). Thus the massless modes of the DD and DN sectors gives an effective D-string world-sheet theory which is precisely that of the heterotic string [6].

Consider now the theory obtained by modding out the weakly coupled IIB
string by $\tilde{\Omega}$, with 32 NS 9-branes. On the perturbative states, $\tilde{\Omega}$ acts as $(-1)^{F_L}$ and so, of the massless fields in the IIB supergravity multiplet, an $N = 1$ supergravity supermultiplet survives, with bosonic sector consisting of the NS-NS fields $g_{MN}, B^1_{MN}, \Phi, A^1_{M...N}$. The oriented fundamental string and NS 5-brane of the IIB theory, coupling to $B^1_{MN}$, also survive, as do the 9-branes coupling to $A^1_{M...N}$.

As $\tilde{\Omega}$ acts as the parity operator on the IIB D-string world-sheet, there are $(0,1)$ strings or ‘D-strings’ of the orientifolded theory that are non-oriented and can be closed or open. The open D-strings can be thought of as ending on the NS 9-branes, and these give them $SO(32)$ Chan-Paton factors. They are not BPS states and have interactions through which they can break and join. In the strong coupling limit, these D-strings become the fundamental open and closed strings of the dual type I theory, while in the weakly coupled theory, they have tensions of order $1/g$ and their effective dynamics is governed by a Born-Infeld action (without Wess-Zumino term). These D-strings can end on the fundamental string and the solitonic 5-brane. There are then DD D-strings which have both ends lying on the fundamental string or solitonic 5-brane, and DN D-strings, one end of which end of which is free and carries an $SO(32)$ charge, and the other end of which ends on the fundamental string or solitonic 5-brane. There are thus $SO(32)$ charges tethered to the fundamental string or solitonic 5-brane by open D-strings.

In the zero coupling limit, the D-string tension becomes infinite, and these D-strings collapse to zero length. Then the $SO(32)$ charges tethered to the fundamental string or solitonic 5-brane are pulled onto the world-volume, giving rise to an $SO(32)$ current density on the world-volume. As the D-strings collapse, all that survives is the zero-slope limit or massless sector of the D-string spectrum, and this must be the same as the massless sector of the excitations of type I strings ending on a D1 or D5-brane in the weakly-coupled type I theory, as the massless sector is in a short multiplet protected by supersymmetry that can be extrapolated from weak to strong coupling. This means that for the D-strings ending on the fundamental string, all that survives are 8 scalars and 8 right-handed Majorana-Weyl world-sheet fermions from the DD sector and 32 left-handed Majorana-Weyl
world-sheet fermions transforming as a 32 of $SO(32)$ from the DN sector. This gives precisely the right world-sheet structure for the free fundamental heterotic string. What is novel is the picture of what happens to the heterotic string if the coupling constant is finite. Then the $SO(32)$ charges, which are confined to the world-sheet at zero coupling, can move off the world-sheet, but are tethered to it by strings of tension $1/g$. However, these can also break, so that a DN string splits into a DN string and an NN string, which is an open D-string that can then escape from the fundamental string.

It has been realised for some time that the 5-brane in the $SO(32)$ theory should carry $SO(32)$ currents [16-23], and in [24] a world-volume structure was proposed, which was shown to lead to cancellation of all anomalies in [25]. For the D5-brane of the type I theory, the massless world-volume fields arise from type I strings with DD and DN boundary conditions [24]. For $N$ coincident D5-branes, the following 6-dimensional (1,0) supermultiplets with $SO(4)$ R-symmetry arise. The DD sector gives a Yang-Mills multiplet with $Sp(N)$ gauge symmetry and a scalar multiplet (transforming as an $N(2N-1) - 1$ antisymmetric tensor representation of $Sp(N)$, plus a singlet) with 4 scalars $X$ and two spinors transforming as a vector and chiral spinor of $SO(4)$, respectively; the $4X^i$ are the collective coordinates for the brane. The DN sector gives a pseudo-real hypermultiplet with 4 scalars which are a vector of $SO(4)$ and a 4-component fermion which is an $SO(4)$ singlet. These hypermultiplets transform as a $(32, d_a)$ of $SO(32) \times Sp(N)$ where the representation of dimension $d_a = N(2N + 1)$ is the adjoint of $Sp(N)$.

The same massless modes should be present on the heterotic 5-brane world-volume, as the massless sector can be extrapolated to strong type I (weak heterotic) coupling. In the construction of the heterotic string by gauging $\tilde{\Omega}$, the world-sheet structure emerges from D-strings ending on the heterotic 5-brane, and in the weak coupling limit of the heterotic string, the D-strings collapse to zero length and only the massless sector survives. The $SO(32)$ currents on the 5-brane world-sheet arise from $SO(32)$ charges tethered to the 5-brane that are ‘pulled in’ in the weak coupling limit, while the $Sp(2)$ currents arise from non-oriented D-strings with
both ends attached to the 5-brane.

The presence of the 9-branes plays a vital role in the construction and are responsible for the gauge structure. If no 9-branes are added, then the perturbative IIB string can be orbifolded by \((-1)^F_L\) to give the IIA string [26], and no gauge sector emerges. It is not clear whether this can be extended to the full non-perturbative theory, and will be discussed further in [15].

One can apply similar considerations to other string theories. Consider the IIB theory. For weak coupling, the D-string dynamics is governed by fundamental strings ending on the D-string, which is useful as it allows detailed calculation. The D-string can be thought of as being surrounded by a cloud of fundamental strings ending on it, and closed fundamental strings can break off from the tethered strings and propagate into the bulk, mediating the interaction with the bulk degrees of freedom. By duality, at strong coupling the \((1,0)\) string dynamics is governed by \((0,1)\) strings ending on the \((1,0)\) string, and this again can be studied using (the S-dual) conformal field theory. Back at weak IIB coupling, D-strings can end on fundamental strings, but this is not so useful in giving the dynamics of the fundamental strings, as we only have a formulation of the D-string dynamics in terms of the fundamental strings themselves (from which an effective Born-Infeld action can be derived). However, the excitations of the D-string that are in short multiplets can be understood, since they can be extrapolated to strong coupling where they are the short-multiplet states of the fundamental string. In particular, the massless sector of the the D-string is the same as that of the fundamental string, and consists of 8 world-sheet scalars \(X^i\) and 16 world-sheet fermions, which transform as a vector and complex chiral spinor of the transverse \(SO(8)\), respectively. At zero IIB string coupling, the tension of the D-strings becomes infinite and only the massless sector survives, so that the familiar world-sheet degrees of freedom of the fundamental IIB string arise from the D-strings attached to the fundamental string that collapse to points, giving local fields on the fundamental string world-sheet. The picture suggests that at finite coupling, the world-sheet is replaced by a cloud of D-strings and the non-BPS states might be understood
in terms of D-string excitations; it would of course be very interesting to find a non-trivial check of this.

The description of the heterotic string proposed here arises from this picture of the type IIB string on gauging $\tilde{\Omega}$. The D-strings become non-oriented and non-BPS, but the world-sheet structure of the free heterotic string again emerges from D-strings attached to the world-sheet, in the limit in which they collapse to points.

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