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Duality Beyond Global Symmetries: The Fate of the $B_{\mu\nu}$ Field

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Duality between the 'axion' field a and the antisymmetric tensor field $B_{\mu\nu}$ is traced after a nonperturbative effect, gaugino condensation, breaks the Peccei-Quinn (PQ) symmetry $a \to a + c$. Even though the PQ symmetry was at its origin, duality is nevertheless *not* broken by this effect. Below condensation scale, the axion simply gets a mass, but in the 'stringy' version, the $B_{\mu\nu}$ field disappears from the propagating spectrum. Its place is taken by a massive 3-index antisymmetric field $H_{\mu\nu\rho}$ which is the one dual to the massive axion. This is a particular case of a general duality in D-dimensions among *massive* p and $D-p-1$-index antisymmetric tensor fields.

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We know that the existence of global symmetries is at the origin of most duality transformations known so far. Actually, a definite prescription has been used to derive those dualities which consists in (i) gauging the global symmetry and (ii) imposing a constraint that guarantees the gauge field to be a pure gauge. By following this prescription it is simple to see duality among 'massless' p and $D-p-2$ antisymmetric tensor fields in D dimensions, abelian and non-abelian dualities in $2D$ as well as abelian and non-abelian bosonization. However, contrary to what has been sometimes claimed in the literature, the existence of a global symmetry is only a sufficient but not necessary condition for duality. Some previous examples are given in [1]. For a recent discussion on duality without the need of isometries see [2].

Here we report on the resolution of the following puzzle: perturbative $4D$ string theory has in its spectrum a two-index antisymmetric tensor field $B_{\mu\nu}$. Because it only has derivative couplings, $B_{\mu\nu}$ is dual to a pseudoscalar field, the axion a. We can transform back and forth from the $B_{\mu\nu}$ and a formulations as long as the corresponding shift symmetries are preserved. It is known that nonperturbative effects break the PQ symmetry of a giving it a mass, then the puzzle is: what happens to the stringy $B_{\mu\nu}$ field in the presence of non-perturbative effects? Is the duality symmetry also broken by those effects? Is it then correct to forget about the $B_{\mu\nu}$ field, as it is usually done, and work only with a? (Since, unlike the axion, $B_{\mu\nu}$ is the field created by string vertex operators). The answer to these questions is very interesting: duality symmetry is *not* broken by the nonperturbative effects but the $B_{\mu\nu}$ field disappears from the propagating spectrum! Its place is taken by a massive 3-index antisymmetric tensor field $H_{\mu\nu\rho}$ dual to the massive axion.

A detailed description of the process is given in [3]. Here I will just sketch the main steps of the derivation. In $4D$ strings, the antisymmetric tensor belongs to a linear superfield L ($\overline{\mathcal{D}\mathcal{D}}L = 0$), together with the dilaton and the dilatino. For simplicity we only consider the couplings of this field to gauge superfields in global supersymmetry, the most general action is then the D-term of an arbitrary function Φ, $\mathcal{L}_L = \left(\Phi(\hat{L})\right)_D$, with $\hat{L} \equiv L - \Omega$ and Ω the Chern Simons superfield, satisfying $\overline{\mathcal{D}\mathcal{D}}\Omega = W_\alpha W^\alpha$, W_α is the gauge field strength superfield. The duality transformation is obtained by starting with the first order system coupled to an external current J:

$$\exp\{i\mathcal{W}(J)\} = \int DA\, DS\, DV\, \exp\left\{i \int d^4x\ \left(\mathcal{L}(V,S) + (JW_\alpha W^\alpha)_F\right)\right\}$$

Where A is the gauge superfield, V an arbitrary vector super field with the lagrangian $\mathcal{L}(V,S) = (\Phi(V))_D + \left(S\overline{\mathcal{D}\mathcal{D}}(V+\Omega)\right)_F$, and S (the same S of S-duality!) starting life as a Lagrange multiplier chiral superfield.

Integrating out S, implies $\overline{\mathcal{D}\mathcal{D}}(V+\Omega) = 0$ or $V = L - \Omega \equiv \hat{L}$, giving back the original theory. On the other hand integrating first V gives the dual theory in terms of S and A. This is the situation above the condensation scale. Below condensation, however, we have to integrate first the gauge fields, after that we have the same two options for getting the two dual theories, the difference now is that the integration over A breaks the PQ symmetry (if there are at least two condensing gauge groups) and we are left with a duality without global symmetries.

To see this, we will concentrate on the $2PI$ effective action $\Gamma(U,V,S)$ obtained in the standard way for $U \equiv \langle\ \mathrm{Tr} W_\alpha W^\alpha\rangle$ [4]. The important result is that since $\mathcal{W}$ depends on S and J only through the combination $S+J$, we can see that $\delta\Gamma/\delta S = \delta\mathcal{W}/\delta S = \delta\mathcal{W}/\delta J = U$ so $\Gamma(U,S,V) = US + \Xi(U,V)$, where $\Xi(U,V)$ is arbitrary, therefore S appears only linearly in the path integral and its integration gives again a δ-function, but imposing now $\overline{\mathcal{D}\mathcal{D}}V = -U$ instead of the constraint $\overline{\mathcal{D}\mathcal{D}}(V+\Omega) = 0$ above condensation scale. We can then see that there is no linear multiplet implied by this new constraint. This is an indication that the $B_{\mu\nu}$ field is no longer in the spectrum.

The new propagating bosonic degrees of freedom in V are, a scalar component, the dilaton, becoming massive after gaugino condensation and a vector field v^μ dual to a, the pseudoscalar component of S. Instead of showing the details of this duality in components, I will describe the following slightly simplified toy model which has all the relevant properties:

$$\mathcal{L}_{v^\mu,a} = -\frac{1}{2}v^\mu v_\mu - a\partial_\mu v^\mu - m^2a^2$$

If we solve for v^μ we obtain $v_\mu = -\partial_\mu a$, substituting back we find

$$\mathcal{L}_a = \frac{1}{2}\partial^\mu a\partial_\mu a - m^2a^2$$

describing the massive scalar a. On the other hand, solving for a we get $a = -\frac{1}{2m^2}(\partial_\mu v^\mu)$ which gives

$$\mathcal{L}'_{v^\mu} = -\frac{1}{2}v^\mu v_\mu + \frac{1}{4m^2}(\partial_\mu v^\mu)^2.$$

The lagrangian $\mathcal{L}'_{v^\mu}$ also describes a massive scalar given by the longitudinal, spin zero, component of v^μ. We can see that the only component that has time derivatives is v^0, so the other three are auxiliary fields. Furthermore, we can easily compute the propagator for $\mathcal{L}'_{v^\mu}$ giving $\delta_{\mu\nu} - \frac{k_\mu k_\nu}{m^2+k^2}$ which is *identical* to the one obtained recently in [5] in their discussion of the axion mass. Therefore we are providing a lagrangian description of that process in terms of a vector field *à la* Duffin-Kemmer or, equivalently, a massive 3-index antisymmetric tensor field. Notice that for $m = 0$, we recover the standard duality among a massless axion and $B_{\mu\nu}$ field. Therefore, after the gaugino condensation process, the original $B_{\mu\nu}$ field of the linear multiplet is projected out of the spectrum in favour of a massive scalar field corresponding to the longitudinal component of v^μ or to the transverse component of the antisymmetric tensor $H_{\mu\nu\rho} \equiv \epsilon_{\mu\nu\rho\sigma} v^\sigma$. Thus solving the puzzle of the axion mass in the two dual formulations.

We can easily see that this duality among *massive* fields is only the $4D$ version of a general duality in D-dimensions among massive p and $D-p-1$-index antisymmetric tensor fields carrying $(D-1)!/p!(D-p-1)!$ propagating degrees of freedom as it can be seen starting from the first order lagrangian

$$\mathcal{L} = \left(H_{M_1\cdots M_p}\right)^2 - \Lambda_{M_1\cdots M_{D-p-1}}\left(\partial_{M_{D-p}} H_{M_{D-p+1}\cdots M_D}\right) - M^2\left(\Lambda_{M_1\cdots M_{D-p-1}}\right)^2$$

and perform the duality transformation as before. Contrary to the massless case, the selfdual models are odd dimensional $D = 2p+1$ [1].

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