Superspace gauge-invariant formulation of a massive tridimensional 2-form field

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By dimensional reduction of a massive supersymmetric $B\wedge F$ theory, a manifestly $N=1$ supersymmetric completion of a massive antisymmetric tensor gauge theory is constructed in $2+1$ dimensions. In $N=1-D=3$ superspace, a new topological term is used to give mass to the Kalb-Ramond field. We introduce a massive gauge invariant model using the Stückelberg formalism and an Abelian topologically massive theory for the Kalb-Ramond superfield. An equivalence of both massive models is suggested. Further, a component field analysis is performed, showing a second supersymmetry in the model.

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I. INTRODUCTION

Antisymmetric tensor fields appear in many field theories. In particular, the Kalb-Ramond (KR) gauge field plays an important role in strong-weak coupling dualities among string theories [1] and in axionic cosmic strings [2]. On the other hand, a first order formulation of the non-Abelian Yang-Mills gauge theory (BF YM model) [3,4] makes use of a two form gauge potential B to contribute to a discussion of the problem of quark confinement in continuum QCD [5]. Another interesting aspect of the $(3+1)$-dimensional $B\wedge F$ term ($F=dA$ is the field strength of a one form gauge potential A) is its ability to give rise to gauge invariant mass to the gauge field [6]. This property has been used to obtain an axion field topologically massive and an axionic charge on a black hole as well [7]. In addition, the existence of the Higgs mechanism to the Kalb-Ramond gauge fields was demonstrated by Rey [8] in the context of closed strings. On the other hand, if coupled to open strings, the KR field becomes a massive vector field through the Stückelberg mechanism. Also, we can mention a topologically massive Kalb-Ramond field in a $D=3$ context that was introduced in Ref. [9].

It is known that massless string excitations may be described by a low-energy supergravity theory and that a massless gravity supermultiplet of graviton, dilaton and Kalb-Ramond fields appears in all known string theories. However, the spectrum of the $D=4$ [10] and $D=3$ [11] compactified theory from $D=10$ supergravity, contains the massive antisymmetric tensor fields. Thus, since supersymmetry places severe constraints on the ground state and the mass spectrum of the excitations, supersymmetric mechanisms of mass generation are of considerable importance. In particular, alternative methods for mass generation can play an important role in the context of superstrings.

In this work we review some aspects of the $N=1-D=4$ BF model. In particular we call attention about a fermionic topological term (to the best of our knowledge, this term has not been discussed in the literature). However, the main new results of this paper are the construction in a $N=1$ superspace of a $N=2-D=3$ topological model, and the $N=1-D=3$ superspace mass generation mechanism for the Kalb-Ramond field. The latter is constructed using a term that we call $B\varphi$, obtained from a dimensional reduction of the $N=1-D=4$ BF model.

This work is organized as follows. In Sec. II we construct an $N=1-D=4$ superspace version of the $U(1)$ BF model. In Sec. III, by means of a dimensional reduction procedure, we obtain a massive antisymmetric tensor field into a $N=2-D=3$ supersymmetric topological massive gauge invariant theory. In contrast to several works on $D=3$ BF models, we have considered here a topological term which involves a KB and a pseudoscalar field with derivative coupling. In Sec. IV an alternative model with an explicit mass breaking term is constructed in $N=1$ superspace and a supersymmetric version of the Stückelberg transformation [12] is used to restore the gauge invariance of the model. In Sec. V we have addressed a $N=1$ superspace mechanism to generate mass for Kalb-Ramond field without loss of gauge invariance. Actually, this mechanism is a superspace version of the topological massive formulation of Deser, Jackiw, and Templeton [13]. Finally, our results are summarized in Sec. VI.

II. THE $N=1-D=4$ EXTENDED BF MODEL

Let us begin by introducing the $N=1-D=4$ supersymmetric BF extended model. For extended we mean that we include mass terms for the Kalb-Ramond field. This mass term will be introduced here for later comparison to the tridimensional case. Actually, this construction can be seen as a superspace and Abelian version of the so-called BF-Yang-Mills models [3].

As our basic superfield action we take[1]

$$S_{BF}^{SS}=\frac{1}{8}\int d^4x\left\{-i\kappa\left[\int d^2\theta B^{\alpha}W_{\alpha}-\int d^2\overline{\theta}\overline{B}_{\dot\alpha}\overline{W}^{\dot\alpha}\right]+\frac{g^2}{2}\left[\int d^2\theta B^{\alpha}B_{\alpha}+\int d^2\overline{\theta}\overline{B}_{\dot\alpha}\overline{B}^{\dot\alpha}\right]\right\}, \tag{2.1}$$

where W_α is a spinor superfield strength, B_α is a chiral spinor superfield, $\overline{D}_{\dot\alpha}B_\alpha=0$, κ and g are massive parameters. Their corresponding θ expansions are

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[1]Our spinorial notations and other conventions follow Ref. [14].

$$W_\alpha(x,\theta,\bar{\theta})=4i\lambda_\alpha(x)-[4\delta_\alpha^\beta D(x)+2i(\sigma^\mu\bar{\sigma}^\nu)_\alpha^\beta F_{\mu\nu}(x)]\theta_\beta+4\theta^2\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu\bar{\lambda}^{\dot{\alpha}} \tag{2.2}$$

$$B_\alpha(x,\theta,\bar{\theta})=e^{i\theta\sigma^\mu\bar{\theta}\partial_\mu}[i\psi_\alpha(x)+\theta^\beta T_{\alpha\beta}(x)+\theta\theta\xi_\alpha(x)], \tag{2.3}$$

where

$$T_{\alpha\beta}=T_{(\alpha\beta)}+T_{[\alpha\beta]}=-4i(\sigma^{\mu\nu})_{\alpha\beta}B_{\mu\nu}+2\varepsilon_{\alpha\beta}(M+iN). \tag{2.4}$$

Our conventions for supersymmetric covariant derivatives are

$$D_\alpha\equiv\frac{\partial}{\partial\theta^\alpha}+i\sigma^\mu_{\alpha\dot{\alpha}}\bar{\theta}^{\dot{\alpha}}\partial_\mu,$$

$$\bar{D}_{\dot{\alpha}}\equiv-\frac{\partial}{\partial\bar{\theta}^{\dot{\alpha}}}-i\theta^\alpha\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu. \tag{2.5}$$

We call attention to the electromagnetic field strength and the antisymmetric gauge field which are contained in W_α and B_α, respectively. In terms of the components fields, the action (2.1) can be read as

$$\begin{aligned}S=\int d^4x\Bigg\{\Bigg[&-\frac{i\kappa}{2}[\xi\lambda-\overline{\xi\lambda}]\\&+\frac{\kappa}{2}[\psi^\alpha\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu\bar{\lambda}^{\dot{\alpha}}+\bar{\psi}_{\dot{\alpha}}(\bar{\sigma}^\mu)^{\dot{\alpha}\alpha}\partial_\mu\lambda_\alpha]\\&+\frac{\kappa}{2}B^{\mu\nu}\tilde{F}_{\mu\nu}-\kappa DN\Bigg]\\&+g^2\Bigg[\tfrac{1}{8}(\psi\xi+\overline{\psi\xi})+\frac{1}{2}B^{\mu\nu}B_{\mu\nu}-\frac{1}{2}(M^2+N^2)\Bigg]\Bigg\}\\=\int d^4x\Bigg[&\left(\frac{i\kappa}{2}\bar{\Xi}\gamma^5\Lambda+\frac{\kappa}{2}\bar{\Psi}\gamma^\mu\partial_\mu\Lambda+\frac{\kappa}{2}B^{\mu\nu}\tilde{F}_{\mu\nu}-\kappa DN\right)\\&+g^2\left(\tfrac{1}{8}\bar{\Psi}\Xi+\tfrac{1}{2}B^{\mu\nu}B_{\mu\nu}-\tfrac{1}{2}(M^2+N^2)\right)\Bigg].\end{aligned} \tag{2.6}$$

In the last equality above, the fermionic fields have been organized as four-component Majorana spinors as follows

$$\Xi=\begin{pmatrix}\xi_\alpha\\ \bar{\xi}^{\dot{\alpha}}\end{pmatrix};\quad\Lambda=\begin{pmatrix}\lambda_\alpha\\ \bar{\lambda}^{\dot{\alpha}}\end{pmatrix};\quad\Psi=\begin{pmatrix}\psi_\alpha\\ \bar{\psi}^{\dot{\alpha}}\end{pmatrix}, \tag{2.7}$$

and we denote the dual field-strength defining $\tilde{F}_{\mu\nu}\equiv\frac{1}{2}\varepsilon_{\mu\nu\alpha\beta}F^{\alpha\beta}$. Furthermore, we use the following identities

$$\begin{aligned}\bar{\Psi}\Lambda&=\overline{\psi\lambda}+\psi\lambda,\\ \bar{\Psi}\gamma^5\Lambda&=\overline{\psi\lambda}-\psi\lambda,\\ \bar{\Psi}\gamma^\mu\Lambda&=\psi\sigma^\mu\bar{\lambda}+\overline{\psi\sigma^\mu}\lambda.\end{aligned} \tag{2.8}$$

The superfield action (2.1) is a particular case of the action proposed in Ref. [15]. However, a point of difference must be noted. In contrast with [15], we have not considered coupling with matter fields and a propagation term for the gauge fields. On the other hand, our superspace *BF* term was constructed in a distinct and simpler way. A quite similar construction was introduced by Clark *et al.* [16].

The off-diagonal mass term $\xi\lambda$ (or $\bar{\Xi}\gamma^5\Lambda$) has been shown by Brooks and Gates [17] in the context of super-Yang-Mills theory. Note that the identity

$$\gamma_5\sigma^{\mu\nu}=\frac{i}{2}\varepsilon_{\mu\nu\alpha\beta}\sigma^{\alpha\beta} \tag{2.9}$$

reveals a connection between the topological behavior denoted by the Levi-Civita tensor $\varepsilon_{\mu\nu\alpha\beta}$, and the pseudoescalar γ_5.

So, it is worthwhile to mention that this term has topological origin and it can be seen as a fermionic counterpart of the *BF* term. In our opinion, this fermionic mass term deserves more attention and will be investigated elsewhere.

III. THE $N=2-D=3$ TOPOLOGICAL MODEL

As it is well known, the *BF* model in $D=3$ consists in a one form field (''B'' field) and one form gauge field A. So, the Chern-Simons term is simply the identification of B and A. However, as has been shown in Ref. [9], after dimensional reduction of the four-dimensional *BF* model, an interesting additional term arises, namely, a topological term which involves a 2 form and a 0 form. We will call it a $B\varphi$ term. A quite similar model was presented in a Yang-Mills version by Del Cima *et al.* [18], and its finiteness was proved in the framework of algebraic renormalization.

Following the procedure of Ref. [9], we will carry out a dimensional reduction in the bosonic sector of Eq. (2.6). Dimensional reduction is usually done by expanding the fields in normal modes corresponding to the compactified extra dimensions, and integrating out the extra dimensions. This approach is very useful in dual models and superstrings [19]. Here, however, we only consider the fields in higher dimensions to be independent of the extra dimensions. In this case, we assume that our fields are independent of the extra coordinate x_3.

Therefore, after dimensional reduction, the bosonic sector of Eq. (2.6) can be written as

$$S_{bos.}=\int d^3x\{[\kappa\varepsilon_{\mu\alpha\beta}V^\mu F^{\alpha\beta}+\kappa\varepsilon_{\mu\nu\alpha}B^{\mu\nu}\partial^\alpha\varphi-\kappa DN]+g^2[\tfrac{1}{2}B^{\mu\nu}B_{\mu\nu}-V^\mu V_\mu-\tfrac{1}{2}(M^2+N^2)]\}, \tag{3.1}$$

where V^μ is a vectorial field and φ represents a real scalar field. Notice that the first term in right-hand side of Eq. (3.1) can be transformed in the Chern-Simons term if we identify $V^\mu\equiv A^\mu$. The second one is the so-called $B\varphi$ term.

Now let us proceed to the dimensional reduction of the fermionic sector of the model. First, note that the Lorentz group in three dimensions is $SL(2,R)$ rather than $SL(2,C)$ in

$D=4$. Therefore, Weyl spinors with four degrees of freedom will be mapped into Dirac spinors.[2] So the correct associations keeping the degrees of freedom are sketched as

$$\Xi=\begin{pmatrix}\xi_\alpha\\ \bar{\xi}^{\dot\alpha}\end{pmatrix}\rightarrow\Xi_\pm=\xi_\alpha\pm i\tau_\alpha,$$

$$\Lambda=\begin{pmatrix}\lambda_\alpha\\ \bar{\lambda}^{\dot\alpha}\end{pmatrix}\rightarrow\Lambda_\pm=\lambda_\alpha\pm i\rho_\alpha,$$

$$\Psi=\begin{pmatrix}\psi_\alpha\\ \bar{\psi}^{\dot\alpha}\end{pmatrix}\rightarrow\Psi_\pm=\psi_\alpha\pm i\chi_\alpha. \tag{3.2}$$

From Eq. (3.2), we find that

$$\Psi\overline{\Xi}\rightarrow\tfrac{1}{2}(\Psi_+\Xi_-+\Psi_-\Xi_+),$$

$$\overline{\Psi}\gamma^\mu\partial_\mu\Lambda\rightarrow\tfrac{1}{2}(\Psi_+\gamma^{\hat\mu}\partial_{\hat\mu}\Lambda_-+\Psi_-\gamma^{\hat\mu}\partial_{\hat\mu}\Lambda_+),$$

$$\Xi\gamma^5\Lambda\rightarrow\tfrac{1}{2}(\Xi_+\Lambda_++\Xi_-\Lambda_-), \tag{3.3}$$

where *hatted* index means three-dimensional space-time.

Thus, the dimensionally reduced fermionic sector of Eq. (2.6) may be written

$$S_{ferm.}=\int d^3x\left\{\frac{i\kappa}{4}(\Xi_+\Lambda_++\Xi_-\Lambda_-)+\frac{\kappa}{4}(\Psi_+\gamma^{\hat\mu}\partial_{\hat\mu}\Lambda_- +\Psi_-\gamma^{\hat\mu}\partial_{\hat\mu}\Lambda_+)+\frac{g^2}{16}(\Psi_+\Xi_-+\Psi_-\Xi_+)\right\}. \tag{3.4}$$

The action $S=S_{bos.}+S_{ferm.}$ is invariant under the following supersymmetry transformations (here and in rest of the paper, greek indices mean three-dimensional space-time):

$$\delta\lambda_\alpha=-iD\eta_\alpha-(\sigma^\mu\sigma^\nu)_\alpha^\beta\eta_\beta F_{\mu\nu},$$

$$\delta\rho_\alpha=iD\zeta_\alpha-(\sigma^\mu\sigma^\nu)_\alpha^\beta\zeta_\beta F_{\mu\nu},$$

$$\delta F^{\mu\nu}=i\partial^\mu(\eta\sigma^\nu\rho-\lambda\sigma^\nu\zeta)-i\partial^\nu(\eta\sigma^\mu\rho-\lambda\sigma^\mu\zeta),$$

$$\delta D=\partial_\mu(-\eta\sigma^\mu\rho+\lambda\sigma^\mu\zeta), \tag{3.5}$$

$$\delta(\psi_\alpha\pm i\chi_\alpha)=\delta\Psi_\pm=i\eta^\beta\tilde{T}_{\beta\alpha}\pm\zeta^\beta\tilde{T}_{\beta\alpha},$$

$$\delta\tilde{T}_{\beta\alpha}=-\eta_\beta\xi_\alpha+\zeta^\lambda\sigma^\mu_{\beta\lambda}\partial_\mu\psi_\alpha,$$

$$\delta(\xi_\alpha\pm i\tau_\alpha)=\delta\Xi_\pm=-i\zeta_\lambda(\sigma^\mu)^{\lambda\beta}T_{\beta\alpha}\mp\eta_\lambda(\overline{\sigma^\mu})^{\beta\lambda}T_{\beta\alpha}, \tag{3.6}$$

where η and ζ are supersymmetric parameters, which indicates that we have two supersymmetries in the aforementioned action.

[2]For details about spinorial dimensional reduction, we suggest Refs. [20] and [21].

IV. REMARKS ON SOME 3D SUPERSYMMETRIC MODELS AND STÜCKELBERG FORMULATION

From the two topological terms introduced in Eq. (3.1) we can setup two supersymmetric models. The first one, which involves a two and a zero form, can be expressed as

$$S=\int d^3x d^2\theta(D^\alpha\Phi B_\alpha+\tfrac{1}{2}g^2B^\alpha B_\alpha), \tag{4.1}$$

where B_α and Φ are spinor and real scalar superfields, which are defined by projection as

$$B_\alpha|=\chi_\alpha,$$

$$D_{(\beta}B_{\alpha)}|=2iM_{\beta\alpha}=M_{\alpha\beta}=B^{\mu\nu}(\sigma_{\mu\nu})_{\alpha\beta},$$

$$D^\alpha B_\alpha|=2N,$$

$$D^\beta D_\alpha B_\beta|=2\omega_\alpha, \tag{4.2}$$

and

$$\Phi|=\varphi,$$

$$D_\alpha\Phi|=\psi_\alpha,$$

$$D^2\Phi|=F. \tag{4.3}$$

Here the supersymmetry covariant derivative is given by $D_\alpha=\partial_\alpha+i\theta^\beta\partial_{\alpha\beta}$. So, in terms of components fields, the action (4.1) becomes

$$S=\int d^3x[(\kappa\partial^{\alpha\beta}\varphi M_{\beta\alpha}+2\kappa\psi^\alpha\omega_\alpha-2\kappa FN) +\tfrac{1}{2}g^2(4\omega^\alpha\chi_\alpha+2i\chi_\alpha\partial^{\beta\alpha}\chi_\beta+M^{\beta\alpha}M_{\alpha\beta}+2N^2)]. \tag{4.4}$$

Starting from the definitions of two spinor superfields given by

$$\Lambda_\alpha|=\xi_\alpha,$$

$$D_{(\beta}\Lambda_{\alpha)}|=2iV_{\beta\alpha},$$

$$D^\alpha\Lambda_\alpha|=2G,$$

$$D^\beta D_\alpha\Lambda_\beta|=2\rho_\alpha, \tag{4.5}$$

and

$$W_\alpha|=\lambda_\alpha,$$

$$D_\alpha W_\beta|=f_{\alpha\beta}, \tag{4.6}$$

where

$$V_{\beta\alpha}\equiv V^\mu(\tilde{\sigma}_\mu)_{\beta\alpha};\quad f_{\alpha\beta}\equiv(\tilde{\sigma}_\mu)_{\alpha\beta}f^\mu;\quad f^\mu=-\frac{i}{2}\varepsilon^{\mu\nu\rho}F_{\nu\rho}, \tag{4.7}$$

we can propose another supersymmetric action, now involving two 1-forms, namely

$$S=\int d^3x d^2\theta(\Lambda^\alpha W_\alpha - g^2\Lambda^\alpha\Lambda_\alpha)$$
$$=\int d^3x[(2\rho^\alpha\lambda_\alpha - iV^{\alpha\beta}f_{\beta\alpha}) - g^2(4\rho^\alpha\omega_\alpha+2i\xi_\alpha\partial^{\beta\alpha}\xi_\beta+V^{\beta\alpha}V_{\beta\alpha}+2G^2)]. \tag{4.8}$$

It is easy to see that the superspace actions (4.1) and (4.8) are not invariant under the following gauge transformations

$$\delta B^\alpha = D^\beta D^\alpha \Pi_\beta,$$
$$\delta\Phi=0, \tag{4.9}$$
$$\delta\Lambda^\alpha = D^\alpha\Omega,$$
$$\delta W^\alpha = 0. \tag{4.10}$$

However, if we reparametrize Λ^α and B^α through introduction of the Stückelberg superfields[3] Θ and Σ_α such that

$$\Lambda^\alpha\rightarrow(\Lambda^\alpha)'=\Lambda^\alpha+\frac{1}{g}D^\alpha\Theta,$$
$$B^\alpha\rightarrow(B^\alpha)'=B^\alpha+D^\beta D^\alpha\Pi_\beta, \tag{4.11}$$

and imposing that Θ and Σ_α transform like

$$\delta\Theta=-g\Omega,$$
$$\delta\Sigma^\beta=-\Pi^\beta, \tag{4.12}$$

we ensure gauge invariance for that superactions.

We remark that integrating out the superfield B_α in Eq. (4.1) we arrive at a supersymmetric Klein-Gordon action and, if we do the same for Λ_α in Eq. (4.8), we obtain a Maxwell superaction. Observe that both these relations may be understood as two duality transformations. We recall here that an analogous connection in $4D$ pure bosonic *BF*-theory was viewed as a perturbative expansion in the coupling g around the topological pure *BF* theory [4]. Thereupon, it may be interesting to perform a similar investigation in the framework of action (4.1).

V. $N=1$ SUPERSPACE TOPOLOGICAL MASS GENERATION

In order to show the topological mass generation for the Kalb-Ramond two form field, we will construct a variation from the model (4.1), by introducing the propagation term for it. Before that, for illustration purpose, we quote the bosonic action introduced in Ref. [9]:

$$S=\int d^3x[\tfrac{1}{6}H_{\mu\nu\rho}H^{\mu\nu\rho}+k\epsilon_{\mu\nu\rho}B^{\mu\nu}\partial^\rho\phi+\tfrac{1}{2}\partial_\mu\phi\partial^\mu\phi], \tag{5.1}$$

where $H_{\mu\nu\rho}$, a three form field strength of the $B^{\mu\nu}$ field, is defined as

$$H_{\mu\nu\rho}=\partial_{[\mu}B_{\mu\rho]}=\partial_\mu B_{\nu\rho}+\partial_\nu B_{\rho\mu}+\partial_\rho B_{\mu\nu}. \tag{5.2}$$

The $N=1$ superspace construction of the supersymmetric version of Eq. (5.1) proceeds as follows. First, we introduce a scalar superfield G defined by

$$G=-D^\alpha B_\alpha, \tag{5.3}$$

where B_α is the super-Kalb-Ramond field defined in Eq. (4.2). Then, after looking the expression (4.1), we find the action

$$S=\int d^3x d^2\theta[-\tfrac{1}{2}(D^\alpha G^2)+kB^\alpha D_\alpha\Phi-\tfrac{1}{2}D^\alpha\Phi D_\alpha\Phi]. \tag{5.4}$$

Now it is straightforward to show that the topological term $kB^\alpha D_\alpha\Phi$ gives rise to a mass term for the super-Kalb-Ramond field. The equation of motion associated with Φ is

$$D^\alpha(kB_\alpha - D_\alpha\Phi)=0. \tag{5.5}$$

Consequently,

$$kB_\alpha - D_\alpha\Phi=\mathcal{C}. \tag{5.6}$$

Since that the constant $\mathcal{C}$ can be absorbed by B_α, we conclude that

$$kB_\alpha - D_\alpha\Phi=0. \tag{5.7}$$

Therefore the original action (5.4) can be rewritten as

$$S=\int d^3x d^2\theta[(D^\alpha G^2)+\tfrac{1}{2}k^2B^\alpha B_\alpha]. \tag{5.8}$$

This exhibits a topological mechanism of mass generation for the Kalb-Ramond field. Naturally, the topological mass terms arise due to the coupling of the B_α and Φ superfields. In other words, this mass term results of the breakdown of the gauge invariance (4.9).

Incidentally let us mention a possible equivalence similar to that between massive topologically and self-dual theories in $D=3$ [13]. Indeed, starting from Eq. (4.1), we can construct an action by introduction of a mass term for the superfield Φ, namely

$$S=\int d^3x d^2\theta(D^\alpha\Phi B_\alpha+\tfrac{1}{2}g^2B^\alpha B_\alpha+m\Phi^2). \tag{5.9}$$

It is easy to see that the equations of motion of Eqs. (5.9) and (5.4) are equivalent. So, the action (5.9) can be consid-

[3]For historical reasons, it is important to cite here the first work, to the best of our knowledge, in the framework of supersymmetric Stückelberg formalism, namely Ref. [22].

ered locally equivalent to action (5.4). On the other hand, it would be interesting to investigate if this equivalence is preserved at quantum level.

VI. CONCLUSIONS

In this work, we have constructed an $N=1-D=3$ superspace action for a model involving an antisymmetric gauge field. Our main point is a topological term that consists in a coupling of this 2-form field and a scalar field. To the best of our knowledge, in the form presented here, this model is completely new in the literature. A similar approach, but involving a 3-form and a scalar fields in $N=1-D=4$, was introduced in Ref. [23].

Starting from the so-called $B\wedge F$ model in $N=1-D=4$ superspace, we carried out a dimensional reduction to the three-dimensional space-time, in order to obtain our basic model. The superspace construction for the $B\wedge F$ is known, but we point out the appearance of a fermionic counterpart of the $B\wedge F$ term.

We have introduced two massive gauge invariant models for an antisymmetric tensor field into a $N=1-D=3$ superspace. In the first, we resort to the Stückelberg formalism and in the other, we construct an abelian topologically massive theory, and a topologically generated mass for the Kalb-Ramond superfield is exhibited. An equivalence of both massive models is suggested. Furthermore, a component field analysis is performed, showing a second supersymmetry in the model.

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