Dual embedding of the Lorentz-violating electrodynamics and Batalin-Vilkovisky quantization.

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Abstract

Modifications of the electromagnetic Maxwell Lagrangian in four dimensions have been considered by some authors \cite{1}. One may include an explicit massive term (Proca) and a topological but not Lorentz-invariant term within certain observational limits. We find the dual-corresponding gauge invariant version of this theory by using the recently suggested gauge embedding method. We enforce this dualisation procedure by showing that, in many cases, this is actually a constructive method to find a sort of parent action, which manifestly establishes duality. We also use the gauge invariant version of this theory to formulate a Batalin-Vilkovisky quantization and present a detailed discussion on the excitation spectrum.

1 Introduction

It has been considered by Carroll, Field and Jackiw, the possibility of modifying the electromagnetic Maxwell Lagrangian to include, for instance, explicit mass term and a topological term \cite{1},

\begin{equation}
L_\mu [A_\mu] = -\frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{m^2}{2} A_\mu A^\mu - L_{CS}[A_\mu]
\end{equation}

where the \( L_{CS} \) is a 3+1 version of the CS action, which couples the dual electromagnetic tensor to an external four vector \( p \):

\begin{equation}
L_{CS}[A_\mu] = -\frac{1}{4} p_\alpha A_\beta \epsilon^{\alpha\beta\mu\nu} F_{\mu\nu}.
\end{equation}

As it has been analysed, if this vector is fixed to be (say) covariantly constant, \( L_{CS} \) is gauge invariant but not Lorentz invariant. The experimental limits to the variations of Maxwell model are also in ref. \cite{1}.

This paper has a two-fold purpose: to construct a gauge invariant version of the Maxwell modified theory via the gauging iterative Noether method and to perform Batalin-Vilkovisky quantization; and furthermore, to discuss on some properties of this procedure of dualization; for instance, its relation to the parent action approach.
Recently the so-called gauging iterative Noether Dualization Method (NDM) [2] has been shown to be effective in establishing some dualities between models [3]. This method is based on the traditional idea of a local lifting of a global symmetry and may be realized by an iterative embedding Noether counter terms. However, this method provides a strong suggestion of duality since it has been shown to give the expected result in the paradigmatic duality between the so-called Self-Dual model and Maxwell-Chern-Simons in three dimensions (SD-MCS). This well-known correspondence was first established in detail by Deser and Jackiw [4], and may be shown using a parent action approach [5].

In this work we argue in favor of this this method (NDM), since we show that, in certain cases, this may actually be seen as a constructive procedure to find a parent action; which would constitute a manifest proof for the duality.

2 The construction of the gauge invariant Lagrangian.

Let us consider the massive Maxwell-Chern-Simons Lagrangian in four dimensions

\[ L^{(0)} = -\frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{m^2}{2} A_\mu A^\mu - \frac{1}{4} p_\alpha A_\beta \epsilon^{\alpha\beta\mu\nu} F_{\mu\nu}. \] (3)

Next, we are going to localize the gauge symmetry,

\[ A_\alpha \rightarrow A_\alpha + \partial_\alpha \eta, \quad \delta A_\alpha = \partial_\alpha \eta, \] (4)

and calculate the first variation of this lagrangian,

\[ \delta L^{(0)}[A_\mu] = (m^2 A_\mu + \beta (\partial^\nu F_{\nu\mu}) - \epsilon_{\alpha\beta\nu\mu} p^\alpha (\partial^\beta A^\nu)) \partial^\mu \eta. \] (5)

We may recognize the Noether current as

\[ J_\mu = m^2 A_\mu + \beta (\partial^\nu F_{\nu\mu}) - \epsilon_{\alpha\beta\nu\mu} p^\alpha (\partial^\beta A^\nu) \] (6)

so that we construct the first iterated lagrangian by introducing an ancillary field \( B \), \( L^{(1)} = L^{(0)} - JB \). If we admit now that \( B \) transforms according to

\[ \delta B_\mu = \delta A_\mu = \partial_\mu \eta, \] (7)

then

\[ \delta L^{(1)} = -(\delta J_\mu) B^\mu. \] (8)

In the other hand we have

\[ \delta J_\mu = m^2 \delta A_\mu = m^2 (\partial_\mu \eta). \] (9)

Note that by defining the second iterated lagrangian by

\[ L^{(2)} = L^{(1)} + \frac{m^2}{2} B_\mu B^\mu, \] (10)

and using (8) and (9), we get that the total variation vanishes, \( \delta L^{(2)} = 0 \). Let us write down the explicit form of this action:

\[ L^{(2)} = -\frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{m^2}{2} A_\mu A^\mu - \frac{1}{4} p_\alpha A_\beta \epsilon^{\alpha\beta\mu\nu} F_{\mu\nu} - (m^2 A_\mu + \beta (\partial^\nu F_{\nu\mu}) - \epsilon_{\alpha\beta\nu\mu} p^\alpha (\partial^\beta A^\nu)) B^\mu + \frac{m^2}{2} B_\mu B^\mu. \] (11)

Solving for \( B \) we get the equation of motion

\[ -J + m^2 B = 0 \] (12)
Plugging this back into (11), we obtain the remarkable gauge invariant theory:

\[ L = \frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} \epsilon_{\alpha\beta\nu\rho} (\partial^\beta A^\nu) A^\rho - \frac{1}{2m^2} \left( \epsilon_{\alpha\beta\nu\rho} p^\alpha (\partial^\beta A^\nu) \right)^2 - \frac{\beta^2}{2m^2} (\partial_\mu F^{\mu\nu})^2 + \frac{\beta}{m^2} \epsilon_{\alpha\beta\nu\rho} p^\alpha (\partial^\beta A^\nu) (\partial_\rho F^{\mu\nu}) . \]  

We should observe also that the massive exitation has been transformed in topological mass by this procedure.

### 2.1 On the Noether Dualization Method.

Note the Maxwell term remain unaffected under dualization and this could be not considered \((\beta = 0)\). This seems to be a characteristic of gauge invariant terms in the initial action. Notice further that in the particular limit \(\beta = 0\), the action 11 is, in certain sense, a parent action for this duality, this is a frequent behavior of this dualization method. To prove this, we only must note that the last iterated lagrangian (11) is precisely the so-called parent action. In other words, we must show that by varying this with respect, first to \(A\) and second to \(B\), the actions (13) and (3) are recovered.

Varying \(L^2\) with respect to \(B\), by construction \(^1\), equation 13 is obtained. Then if we variate this action with respect to \(A\), we obtain the equation of motion

\[ J(A - B) = 0. \]  

Next, let us define the new variable \(\hat{A} \equiv A - B\); equation (14) corresponds to an action \(L^{(0)}[\hat{A}]\). This completes our argument.

The similarity of this duality with the one in three dimensions between the Self-Dual model (SD) and Maxwell-CS (MCS) must to be remarked. This is not surprising since, if we choice the external vector \(p_\mu\) to coincide with an (espalial) element of cartesian basis of the space-time, and writing the fields in components, one may directly verify that action (with \(k = 0\)) reduces to SD in three dimensions and the action (13) coincides with MCS. Thus, of course, the duality is preserved *.

Notice that this argument can be repeated in each case studied by this Noether method, except when the non-linearity is marked, such as in the case of the duality between Non-Abelian Self-Dual and Yang-Mills Chern Simons theories (in three dimensions). In these cases another proof must be given, a suggestion to fill this gap was given precisely for this problem in ref. [6]

### 3 Batalin-Vilkovisky Quantization.

After the model has been analysed in terms of its particle excitations and its consistency has been ascertained, our goal is to implement a procedure to obtain a gauge fixed Generating Functional of the new gauge invariant Lagrangian. For this reason we will work with the field antifield formalism.

The standard form of field antifield quantization or BV scheme [8] is based on imposing explicit BRST invariance of the vacuum functional, by considering a closed equation that generates the gauge algebra and uses it afterwards to construct the functional generator of theory. An important feature of the standard BV is the anticanonical form of its phase space, where the coordinates has a relationship conjugated with respect to the antibracket operation. Another important aspect of the field antifield formalism is that canonical transformations [8, 9] do not change the form of the antibracket. At classical level the action is constructed with the same Lagrangian that one considered on the equation (13):

\[ L = \frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} \epsilon_{\alpha\beta\nu\rho} (\partial^\beta A^\nu) A^\rho - \frac{1}{2m^2} \left( \epsilon_{\alpha\beta\nu\rho} p^\alpha (\partial^\beta A^\nu) \right)^2 - \frac{\beta^2}{2m^2} (\partial_\mu F^{\mu\nu})^2 + \frac{\beta}{m^2} \epsilon_{\alpha\beta\nu\rho} p^\alpha (\partial^\beta A^\nu) (\partial_\rho F^{\mu\nu}) , \]

\(^1\)In fact, this constitutes the final step of the Gauging Noether method.
where the gauge symmetry is
\[ \delta A_\mu = \partial_\mu \eta \] (15)
and the corresponding BRST symmetry can be read by
\[ sA_\mu = \partial_\mu C \] (16)
\[ sC = b. \] (17)

Then we can deal now with the non minimal solution of Master equation, by including the auxiliary fields \((b, \bar{C})\)
\[ S = \int d^4x \left( L + A_\mu^* \partial^\mu C + \bar{C}^* b \right) \] (18)

Notice that formally, the terms that breaks the Lorentz invariance, more precisely the external vector \(p\), can be treated as an external source in order to perform a BV quantization of the theory. However, since this is not a field of the non minimal sector of theory its nature is BRST invariant. Therefore the path integral is independent of the gauge choice, and its introduction does not introduce subsequent modifications in the BV scheme to the quantization procedure. We are ready now to gauge fix the model. With this aim we will consider the functional fermion,
\[ \Psi = \int d^4x \bar{C} \left( \partial_\mu A_\mu - \frac{1}{2} b \right), \] (19)
the BRST variation of it is written by
\[ s\Psi = \int d^4x \left( b \partial_\mu A_\mu - \bar{C} \partial_\mu \partial_\mu C - \frac{1}{2} b^2 \right). \] (20)
The generator functional can be read with the gauge fixed action by
\[ Z[p, J] = \int DA_\mu DC D\bar{C} Db DA_\mu^* D\bar{C}^* e^{-i(S + s\Psi + J_\mu A^\mu)}, \] (21)
After functional integration over \(A_\mu^*, \bar{C}^*\) and \(b\) we are able to write the final form for the generating functional
\[ Z[p, J] = \int DA_\mu DC e^{-i \int d^4x (L + \partial_\mu b A^\mu + \bar{C} \partial_\mu C + J_\mu A^\mu)}, \] (22)
where \(L\) is the same written on (13).

4 The Spectrum of the Theory and Unitarity Requirements.

We now intend to investigate the pole structure of the gauge propagator stemming from our Lagrangian
\[ L = \frac{\beta}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} \varepsilon_{\alpha\beta\nu\rho} \varepsilon^\rho (\partial^\beta A^\nu) A^\mu - \frac{1}{2m^2} (\varepsilon_{\alpha\beta\nu\rho} \varepsilon^\rho (\partial^\beta A^\nu))^2 \]
\[ - \frac{\beta^2}{2m^2} (\partial_\mu F^{\mu\nu})^2 + \frac{\beta}{m^2} \varepsilon_{\alpha\beta\nu\rho} \varepsilon^\rho (\partial^\beta A^\nu) (\partial_\mu F^{\mu\nu}) - \frac{1}{2\alpha} (\partial_\mu A^\mu), \] (23)
in which we have added a gauge-fixing term. By means of partial integrations in the action, we can rearrange the Lagrangean in terms of spin operators
\[ L = \frac{1}{2} A^\mu \left\{ \left[ -\beta \Box - \frac{1}{m^2} (p^2 - \lambda^2) - \frac{\beta^2 \Box^2}{m^2} \right] \theta_{\mu\nu} + \left( \theta_{\alpha\beta} + \frac{\lambda^2}{m^2} \right) \omega_{\mu\nu} + \left( -1 - \frac{2\beta \Box}{m^2} \right) S_{\mu\nu} \right\} = \frac{1}{2} A^\mu \partial_\mu A^\nu, \] (24)
with $\theta_{\mu\nu} = g_{\mu\nu} - \frac{\partial_{\mu} \partial_{\nu}}{\Box}$ and $\omega_{\mu\nu} = \frac{\partial_{\mu} \partial_{\nu}}{\Box}$ being the transversal and the longitudinal operators, respectively, and

\[
S_{\mu\nu} = \epsilon_{\mu\nu\alpha\beta} p^\alpha \partial^\beta, \tag{25}
\]

\[
A_{\mu\nu} = p_{\mu} p_{\nu} \quad \text{and} \quad \tag{26}
\]

\[
\Sigma_{\mu\nu} = p_{\mu} \partial_{\nu}, \tag{27}
\]

generated by the inclusion of the external vector $p^\mu$. The $\lambda$ is just $\Sigma_{\mu}^\mu = p_{\mu} \partial^\mu$. The algebra of these operators is shown in Table 1:

| $\theta_{\mu\nu}$ | $\omega_{\mu\nu}$ | $S_{\mu\nu}$ | $A_{\mu\nu}$ | $\Sigma_{\mu\nu}$ | $\Sigma_{\mu\nu}$ |
|-------------------|-------------------|-------------|--------------|-------------------|-------------------|
| $\theta_{\mu\nu}$ | $\theta_{\mu\nu}$ | $0$          | $S_{\mu\nu}$ | $A_{\mu\nu} - \frac{\alpha}{\Box} \Sigma_{\mu\nu}$ | $\Sigma_{\mu\nu} - \lambda \omega_{\mu\nu}$ | $0$ |
| $\omega_{\mu\nu}$ | $0$               | $\omega_{\mu\nu}$ | $0$          | $\partial^\lambda \Sigma_{\mu\nu}$ | $\lambda \omega_{\mu\nu}$ | $0$ |
| $S_{\mu\nu}$      | $S_{\mu\nu}$      | $0$          | $0$          | $0$               | $0$               | $0$ |
| $A_{\mu\nu}$      | $A_{\mu\nu} - \frac{\alpha}{\Box} \Sigma_{\mu\nu}$ | $0$          | $\lambda A_{\mu\nu}$ | $\lambda \Sigma_{\mu\nu}$ | $\lambda \Sigma_{\mu\nu}$ | $\lambda A_{\mu\nu}$ |
| $\Sigma_{\mu\nu}$ | $\Sigma_{\mu\nu} - \lambda \omega_{\mu\nu}$ | $\lambda \omega_{\mu\nu}$ | $0$          | $v^2 \Sigma_{\mu\nu}$ | $\lambda A_{\mu\nu}$ | $\lambda \Sigma_{\mu\nu}$ |
| $\Sigma_{\nu\mu}$ | $\Sigma_{\nu\mu}$ | $0$          | $0$          | $v^2 \Sigma_{\mu\nu}$ | $\lambda A_{\mu\nu}$ | $\lambda \Sigma_{\mu\nu}$ |
| $\Sigma_{\nu\mu}$ | $\Sigma_{\nu\mu}$ | $0$          | $0$          | $v^2 \Sigma_{\mu\nu}$ | $\lambda A_{\mu\nu}$ | $\lambda \Sigma_{\mu\nu}$ |

Table 1: Multiplicative Table fulfilled by $\theta, \omega, S, \Lambda$ and $\Sigma$. The products are supposed to obey the order "row times column".

Making use of Table 1, we can invert $O_{\mu\nu}$, which yields the following vector propagator in momenta space:

\[
\langle A_{\mu} A_{\nu} \rangle = \frac{i}{D} \left\{ \left[ -m^2 \left( k^2 (k^2 - m^2) - p^2 k^2 + (p \cdot k)^2 \right) \right] \theta_{\mu\nu} \right.
\]

\[
+ \left[ \frac{-\alpha D}{k^2} - \frac{m^2 (p \cdot k)^2 F}{k^2 (k^2 - m^2)} \right] \omega_{\mu\nu} - i m^2 \left( 2k^2 - m^2 \right) S_{\mu\nu}
\]

\[
- \frac{m^2 F}{k^2 - m^2} \partial_{\mu\nu} + \frac{m^2 (p \cdot k) F}{k^2 (k^2 - m^2)} \left( \Sigma_{\mu\nu} + \Sigma_{\nu\mu} \right) \right\}. \tag{28}
\]

The factor $F$ and the denominator $D$ are given by

\[
F = 3k^2 (k^2 - m^2) + m^4 + p^2 k^2 - (p \cdot k)^2 \tag{29}
\]

and

\[
D = \left[ k^4 + p^2 k^2 - (p \cdot k)^2 \right] \left[ (k^2 - m^2)^2 + p^2 k^2 - (p \cdot k)^2 \right]. \tag{30}
\]

In the above expressions, we have made $\beta = 1$, in order to compare our spectrum with the obtained in the work [7], where the original theory (the Lagrangian given by eq. (1) was analysed. It is interesting to note that our denominator is simply the product of two denominators: the one that comes from the original massive theory (eq.(1)) and the one that comes from the non massive theory (eq.(1) without the Proca term). We have now doubled the number of degrees of freedom, since the new Lagrangian has derivatives of higher order. This will be analysed further.

Here we will analyse situations in which the external vector $p^\mu$ is purely space-like. The time- and light-like terms allows for non physical poles in the theory of eq(1) as it is shown in ref. [7]. Then we take $p^\mu = (0, 0, 0, \mu)$ and $k^\mu = (k_0, 0, 0, k_3)$. The denominator, then, reads

\[
D = [k^2 - m_1^2][k^2 - m_2^2][k^2 - m_1^2][k^2 - m_2^2], \tag{31}
\]

with the poles

\[
k_0^2 = m_1^2 = \frac{1}{2} \left[ 2(m^2 + k_3^2) + \mu^2 + \mu \sqrt{\mu^2 + 4(m^2 + k_3^2)} \right], \tag{32}
\]

\[
k_0^2 = m_2^2 = \frac{1}{2} \left[ 2(m^2 + k_3^2) + \mu^2 - \mu \sqrt{\mu^2 + 4(m^2 + k_3^2)} \right], \tag{33}
\]
Besides these poles, we still have the poles \( k_0^2 = k_3^2 \) \((k^2 = 0)\) and \( k_0^2 = m^2 + k_3^2 \) \((k^2 = m^2)\). The poles \( m_1^2, m_2^2 \) and \( m^2 + k_3^2 \) come from the massive sector, while the other three arise from the massless sector.

We start off with the calculus of the residue matrix for the pole \( k_0^2 = m_1^2 \), for which we find

\[
R_1 = \frac{1}{\sqrt{\mu^2 + 4(m^2 + k_3^2)}} \begin{pmatrix}
0 & 0 & 0 & 0 \\
0 & m_1 & im_1 & 0 \\
0 & -im_1 & m_1 & 0 \\
0 & 0 & 0 & 0
\end{pmatrix},
\]

with the eigenvalues

\[
\begin{align*}
\lambda_1 &= 0 \\
\lambda_2 &= 0 \\
\lambda_3 &= \frac{2|m_1|}{\sqrt{\mu^2 + 4(m^2 + k_3^2)}} > 0 \quad \text{and} \\
\lambda_4 &= 0.
\end{align*}
\]

As it can be seen, this pole is to be associated with only one physical degree of freedom, since we have only one non-null eigenvalue. The same happens with the pole \( k^2 = m_2^2 \), for which the results are completely similar with the exchange of \( m_1 \) for \( m_2 \).

We now study the pole \( k_0^2 = m^2 + k_3^2 \). The associated residue matrix reads as below

\[
R_m = \frac{k^2}{m^2} \begin{pmatrix}
1 & 0 & 0 & \frac{|k_0|}{m^2} \left((m^2 + k_3^2) \right)^{1/2} \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
\frac{|k_0|}{m^2} \left((m^2 + k_3^2) \right)^{1/2} & 0 & 0 & \frac{(m^2 + k_3^2)}{m^2}
\end{pmatrix},
\]

with a unique non-vanishing eigenvalue:

\[
\lambda = \frac{m^2 + 2k_3^2}{m^2} > 0.
\]

These results obtained for the three initial poles, coming from the original theory, that includes a Proca term, are exactly the same obtained in [7], where the Lagrangean eq (1) was studied. They express three degrees of freedom, which one coming from one pole.

The study of the three remaining poles brings about an interesting result: the residues found for the poles coming from the massless factor of \( D \) are simply the residues found for the massless theory up to a minus sign. This suggest that maybe the total propagator can be rearranged so as to be expressed as the difference between the propagators of the massive and the non-massive theories. Indeed, we have:

\[
\langle A_\mu A_\nu \rangle = \langle A_\mu A_\nu \rangle_1 - \langle A_\mu A_\nu \rangle_2 - \frac{1}{m^2} \omega_{\mu\nu},
\]

with

\[
\begin{align*}
\langle A_\mu A_\nu \rangle_1 &= \frac{i}{D_1} \left\{ - (k^2 - m^2) \varphi_{\mu\nu} - \frac{(p \cdot k)^2}{(k^2 - m^2)} \omega_{\mu\nu} \\
&\quad - i S_{\mu\nu} - \frac{k^2}{(k^2 - m^2)} \Lambda_{\mu\nu} + \frac{(p \cdot k)}{(k^2 - m^2)} (\Sigma_{\mu\nu} + \Sigma_{\nu\mu}) \right\}.
\end{align*}
\]
\[ \langle A_\mu A_\nu \rangle_2 = \frac{i}{D_2} \left\{ -k^2 \theta_{\mu\nu} - \frac{(p \cdot k)^2}{k^2} \omega_{\mu\nu} \right. \\
left. -iS_{\mu\nu} - A_{\mu\nu} + \frac{(p \cdot k)}{k^2} (\Sigma_{\mu\nu} + \Sigma_{\nu\mu}) \right\}, \tag{45} \]

where we have made the gauge choice \( \alpha = 0 \). This is a remarkable result. The minus sign in the second propagator guarantees us that the physical particles predicted by the original theory are the same in its dual theory, since the negative eigenvalues obtained in the second part are compatible with negative norm states. This apparently surprising result may actually be traced back to the case of the usual Proca theory. If we proceed to the dualisation of the action for the Proca field, as we have done in Section 1 (without the Lorentz-breaking term), we can show that the \( A_\mu \)-propagator of the dual theory also turns out to have the same relation with the massive and the massless \( A_\mu \)-propagators.

5 A Final Comment.

The gauge invariant Maxwell Theory was modified in such a form that the gauge invariance is broken within certain observational limits. We showed here that this invariance can be restored and this modified version of the electromagnetism results to be equivalent to a novel gauge invariant theory (equation (13)). However the Lorentz invariance remains broken.

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References

[1] S. M. Carrol, George B. Field and Roman Jackiw, Phys. Rev. D 41 (1990) 1231
[2] M. A. Anacleto, A. Ilha, J. R. S. Nascimento, R. F. Ribeiro and C. Wotzasek, Phys. Lett. B504 (2001) 268.
[3] A. Ilha and C. Wotzasek, Nucl.Phys. B604 (2001) 426.
[4] S. Deser and R. Jackiw, Phys. Lett. B 139 (1984) 2366.
[5] For a recent review in the use of the master action in proving duality in diverse areas see: S. E. Hjelmeland, U. Lindström, UIO-PHYS-97-03, May 1997. e-Print Archive: hep-th/9705122
[6] M. Botta Cantcheff, Phys.Lett. B528 (2002) 283-287.
[7] A. P. Baêta Scarpelli, H. Belich, J. L. Boldo and J. A. Helayel-Neto, Aspects of causality and unitarity and comments on vortex-like configurations in an abelian model with a Lorentz-breaking term, hep-th/0204232 (to appear in Phys. Rev. D)
[8] I. A. Batalin and G. A. Vilkovisky Phys. Lett. B 102 (1981) 27
[9] E. M. C. Abreu, N. R. F. Braga and C. F. L. Godinho, Nucl. Phys (1987)
[10] P. K. Townsend, K. Pilch and P. van Nieuwenhuizen, Phys. Lett. B 136 (1984) 38.
[11] E. Fradkin and F. A. Schaposnik, Phys. Lett. B338 (1994) 253; E. Fradkin, F. A. Schaposnik, Phys. Lett.B358 (1994) 253.