Finite BRST Transformations for the Bagger-Lambert-Gustavsson Theory

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Abstract

In this paper we analyse the Bagger-Lambert-Gustavsson (BLG) theory in \( \mathcal{N} = 1 \) superspace. Furthermore, we will construct the BRST transformations for this theory. These BRST transformations will be integrated out to obtain the finite field dependent version of BRST (FFBRST) transformations. We will also analyse the effect of the FFBRST transformations on the effective action. We will thus show that the FFBRST transformations can be used to relate generating functionals of the BLG theory in two different gauges.

1 Introduction

Bagger-Lambert-Gustavsson (BLG) theory is thought to be the dual boundary gauge theory to a 11-dimensional supergravity theory living on \( \text{AdS}_4 \times S_7 \). It is thus a superconformal field theory with \( \mathcal{N} = 8 \) supersymmetry and its gauge group is \( \text{SO}(4) \). This theory has been constructed using Lie 3-algebra \([1, 2, 3, 4, 5]\). It has been analysed in \( \mathcal{N} = 1 \) superspace formalism \([6, 7]\). By complexifying the matter fields, the BLG theory can also be written as a Chern-Simons-matter theory with the gauge group \( SU(2) \times SU(2) \) generated by ordinary Lie algebra. One of the gauge groups is associated with Chern-Simons level, \( k \) and the other, with \( -k \) \([8]\). BLG theory only represents two M2-branes, and it has not been possible to express more
than two M2-branes using the BLG theory. However, inspired by BLG theory Aharony-Bergman-Jafferis-Maldacena (ABJM) theory has been constructed, and this theory represents N M2-branes [10,11]. The gauge symmetry in the ABJM theory is generated by an ordinary Lie algebra and the gauge group of this theory is $U(N) \times U(N)$. This theory only has manifest $\mathcal{N} = 6$ supersymmetry, which is expected to get enhanced to $\mathcal{N} = 8$ supersymmetry from a variety of mechanisms [12]. In this paper we will only analyse the BLG theory using Lie 3-algebra, as it is more difficult to write the FFBRST transformations for the ABJM theory.

It may be noted that as the BLG theory has gauge symmetry, it cannot be quantized without getting rid of these unphysical degrees of freedom. This can be done by fixing a gauge. The gauge fixing condition can be incorporated at a quantum level by adding ghost and gauge fixing terms to the original classical Lagrangian. It is known that for a gauge theory the new effective Lagrangian constructed as the sum of the original classical Lagrangian with the gauge fixing and the ghost terms, is invariant a new set of transformations called the BRST transformations [13,14]. BRST symmetry has also been studied in non-linear gauges [15,16]. In these gauges quadratic ghost interactions are introduced and the effective theory is invariant under a larger algebra called the Nakanishi-Ojima algebra [17]. In fact, BRST symmetry for ABJM theory has also been studied [8]. The BRST symmetry can be used to project out the sub-space of physical states from the total Hilbert space. It has been demonstrated that the nilpotency of these transformations is crucial for the unitarity of the $S$-matrix in the $M$-theory [8]. Even though the BRST symmetry of the ABJM theory has been studied, so far the BRST symmetry of based on a Lie 3-algebra has not been studied. It would be interesting to analyse the BRST symmetry directly based on Lie 3-algebra because this structure has been used to study the action for M5-branes [18]. In this proposal the BLG action with Nambu-Poisson 3-bracket has been identified with the M5-brane action with a large worldvolume three form field. In this paper we analyse the infinitesimal BRST symmetry for the BLG theory.

The infinitesimal global BRST transformations can be integrated out to obtain the FFBRST transformations [19]. Various applications of these FFBRST transformations have been studied [19,20,21,22,23,24,25,26]. A correct prescription for the poles in the gauge field propagators in non-covariant gauges has been derived by connecting effective theories in covariant gauges to the theories in noncovariant gauges by using FFBRST transformation [27]. The divergent energy integrals in the Coulomb gauge have also been regularized by modifying the time like propagator by using FFBRST transformation [28]. The Gribov-Zwanziger theory [29,30], which is free from Gribov copies and plays a crucial role in the non-perturbative infrared regime while it can be neglected in the perturbative ultraviolet regime, has also been related to a theory with Gribov copies i.e. Yang-Mills
theory in Euclidean space through FFBRST transformation \cite{3,1,2}. In this paper we will also analyse the FFBRST for the BLG theory and show how it can be used to relate the BLG theory in two different gauges.

## 2 BLG Theory

In this section, first of all we review the construction of BLG theory in $\mathcal{N} = 1$ superspace. To do that we first start from reviewing the basic properties of a Lie 3-algebra. A Lie 3-algebra $\mathcal{A}$ is a vector space with basis $T^A$, $a = 1, \ldots, \dim \mathcal{A}$, endowed with a trilinear antisymmetric product \cite{9},

\[ [T^A, T^B, T^C] = f^{ABC} D^D. \] (1)

The algebra is accompanied by an inner product, $h^{AB} = Tr(T^A T^B)$, with which indices may be raised and lowered. The structure constants of the algebra are required to be totally anti-symmetric, $f^{ABCD} = f_{[ABCD]}$ and satisfy the fundamental identity,

\[ f^{[ABC} G^{D]EH} = 0. \] (2)

The classical Lagrangian density for the BLG theory in this superspace formalism is given by, $L_c = L_M + L_{CS}$, where $L_{CS}$ is the Lagrangian densities for the Chern-Simons theory and $L_M$ is the Lagrangian density for the matter fields. Now, the non-Abelian Chern-Simons theory on this superspace can now be written as

\[ L_c = \frac{k}{4\pi} \int d^2 \theta \ Tr[f^{ABCD} \Gamma_a^{ab} \Omega_a^{aCD}], \] (3)

where $k$ is an integer and

\[ \Omega_{ABa} = \omega_{AB} - \frac{1}{6} C^{CD,EF}_{AB} \Gamma_{b}^{CD} \Gamma_{ab}^{EF} \] (4)

\[ \omega_{ABa} = \frac{1}{2} D_a \Gamma_{ABb} - \frac{i}{2} C^{CD,EF}_{AB} \Gamma_{CD} \Gamma_{b}^{EF} \Gamma_{a}^{aEF} - \frac{1}{6} C^{CD,EF}_{AB} C^{LM,NP}_{EF} \Gamma_{b}^{CD} \Gamma_{b}^{LM} \Gamma_{a}^{NP}, \] (5)

\[ \Gamma_{ABab} = -\frac{i}{2} \left[ D_{a} \Gamma_{ABb} - i C^{CD,EF}_{AB} \Gamma_{aCD} \Gamma_{b}^{EF} \right], \] (6)

here $D_a$ is given by $D_a = \partial_a + (\gamma^\mu \partial_\mu)^b_a \theta_b$. The Lagrangian density for the matter fields is given by

\[ L_M = \frac{1}{4} \int d^2 \theta \ Tr \left[ \nabla^a X^I \nabla_a X_I + V \right], \] (7)

\[ \nabla_a X_I = \partial_a X_I + \frac{i}{2} \left( \gamma^\mu \partial_\mu \right)^b_a X_I \theta_b. \]
where the covariant derivatives are given by $\nabla_a X^{AI} = D_a X^{AI} + i\Gamma^A_{aB} X^B$, and the potential term given by $V = f^A_{BCDF} \epsilon^{KL} [ X^A X^B \tilde{X}^C Y^D ]$. The infinitesimal gauge transformations for these fields are written as,

$$\begin{align*}
\delta X^{IA} &= i(\Lambda X^I)^A, \\
\delta X^{IA}^\dagger &= -i(\Lambda X^I)^A, \\
\delta \Gamma^A_a &= (\nabla_a \Lambda)^A.
\end{align*}$$

The Lagrangian for the BLG theory is invariant under these gauge transformations, $\delta L_{BLG} = 0$, where $\delta L_{BLG} = \delta L_{kcs}(\Gamma) - \delta \tilde{L}_{kcs}(\tilde{\Gamma}) + \delta L_M$. All the degree’s of freedom in the Lagrangian density for this BLG theory are not physical because it is invariant under gauge transformations. So, we have to fix a gauge before doing any calculations. This can be done by choosing the following gauge fixing conditions, $G = 0$, where $G = D^a \Gamma_a$. These gauge fixing conditions can be incorporate at the quantum level by adding the following gauge fixing term to the original Lagrangian density,

$$L_{gf} = \int d^2 \theta \left[ f^A_{BCDF} b_{AB} D^a \Gamma_{aCD} + \frac{\alpha}{2} f^A_{BCDF} b_{AB} b_{CD} \right].$$

The ghost term corresponding to this gauge fixing term can be written as

$$L_{gh} = \int d^2 \theta \left[ f^A_{BCDF} \tau_{AB} D^a \nabla_a c_{CD} \right].$$

3 BRST Symmetry

The total Lagrangian density obtained by addition of the original classical Lagrangian density, the gauge fixing term and the ghost term can be used to construct the effective action for the BLG theory as,

$$S_{BLG} = \int d^3 x [ L_c + L_{gf} + L_{gh}].$$

This effective action is used to define the generating functional for the BLG theory as

$$Z = \int D\Gamma Dc D\tilde{c} D\tau e^{iS_{BLG}}.$$

The total Lagrangian density given in Eq. (11) is invariant under the following infinitesimal BRST transformations,

$$\begin{align*}
\delta \Gamma^A_a &= (\nabla_a \Lambda)^A, \\
\delta c_{AB} &= -\frac{1}{2} C^A_{EF} C^A_{CD} c_{EF}, \\
\delta \tau_{AB} &= b_{AB}, \\
\delta b_{AB} &= 0, \\
\delta X^{IA} &= i c_{AB} X^B, \\
\delta X^{IA}^\dagger &= -i X^B c_{AB}.
\end{align*}$$

where BRST parameter is global, infinitesimal and anticommuting in nature.
Now, to check the nilpotency of such transformations we have

\[ s^2 b^{AB} = 0, \]
\[ s^2 x^{AB} = s b^{AB} = 0, \]
\[ s^2 c_{AB} = -s \frac{1}{2} C^{CD,E,F} c_{CD} c_{E,F} \]
\[ = \frac{1}{4} C_{AB} c^{LM,PT} c_{LM} c_{PT} c_{EF} \]
\[ - \frac{1}{4} C_{AB} c_{LM,PT} c_{CD} c_{LM} c_{PT} = 0, \]
\[ s^2 \Gamma^A = s D_a c_{AB} \]
\[ = D_a [c_{AB} + C^{CD,E,F} \Gamma_{CD} c_{E,F}] \]
\[ + C^{CD,E,F} \Gamma_{CD} D_a c_{E,F} \]
\[ - C^{CD,E,F} \Gamma_{CD} D_a c_{E,F} \]
\[ + C^{CD,E,F} \Gamma_{CD} c_{E,F} \Gamma_{CD} a c_{PQ} c_{RS} = 0, \]
\[ s^2 X_A^I = i s c_{AB} X_B^I \]
\[ = \frac{i}{2} C_{AB} c_{CD} c_{E,F} X_B^I \]
\[ - i c_{AB} c_{BE} X_E^I = 0, \]
\[ s^2 X_A^{I\dagger} = - i s X_B^{I\dagger} c_{AB} \]
\[ = \frac{i}{2} C_{AB} c_{CD} c_{BE} X_B^{I\dagger} \]
\[ + i X_E^{I\dagger} c_{BE} c_{AB} = 0, \]
\[ \text{(14)} \]

Thus, these BRST transformations are nilpotent, \( s^2 = 0 \).

We can now express the sum of the gauge fixing term and the ghost term as

\[ L_{gf} + L_{gh} = \int d^2 \theta \ s \left[ f^{ABCD} c_{AB} \left( D^a \Gamma_{aCD} - \frac{\alpha}{2} b_{CD} \right) \right]. \]
\[ \text{(15)} \]

In fact, the invariance of the total Lagrangian density follows from the nilpotency of the BRST transformations. This is because \( L_{gh} + L_{gf} \) can be expressed as a total BRST variation and hence the action of \( s \) on \( L_{gh} + L_{gf} \) vanishes. The BRST variation of the original theory is the gauge transformation with gauge fields replaced by ghosts or anti-ghosts, respectively,

\[ s L_e + s L_{gh} + s L_{gf} = 0. \]
\[ \text{(16)} \]

### 4 FFBRST Transformation

In this section, we construct the FFBRST transformations for the BLG theory. In order to do that we first define \( \Phi_i(x, \kappa) = \Phi_i^{AB}(x, \kappa) T_A T_B \), where
\( \Phi_{iAB} = (\Gamma^A_B, c^{AB}, \bar{\nu}^{AB}, b^{AB}) \), here all the fields depend on some parameter, \( \kappa : 0 \leq \kappa \leq 1 \), in such a manner that \( \Phi^i(x, 0) \) are the initial fields and \( \Phi^i(x, 1) \) are the transformed field. Now, we also define \( \Theta[\Phi] \) as a functional with odd Grassmann parity. This can obtained from a infinitesimal field dependent parameter through the following relation

\[
\Theta[\Phi(x)] = \epsilon[\Phi(x)] \exp \frac{F[\Phi(x)] - 1}{F'[\Phi(x)]},
\]

where

\[
F = \frac{\delta \epsilon[\Phi(x)]}{\delta \Gamma_a(x)} s\Gamma_a(x) + \frac{\delta \epsilon[\Phi(x)]}{\delta c(x)} sc(x) + \frac{\delta \epsilon[\Phi(x)]}{\delta \bar{\nu}(x)} s\bar{\nu}(x) + \frac{\delta \epsilon[\Phi(x)]}{\delta b(x)} sb(x).
\]  

Now, the infinitesimal parameter in the BRST transformation is made field dependent and hence the BRST transformation can be written as

\[
\frac{d}{d\kappa} \Phi^i(x, \kappa) = s\Phi^i(x, \kappa) \epsilon[\Phi(x), \kappa],
\]

where \( \epsilon[\Phi(x), \kappa] \) is an infinitesimal field dependent parameter. By integrating these equations from \( \kappa = 0 \) to \( \kappa = 1 \), it has been shown that the \( \Phi^i(x, 1) \) are related to \( \Phi^i(x, 0) \) by the FFBRST transformation as

\[
\Phi^i(x, 1) = \Phi^i(x, 0) + s\Phi^i(x, 0)\Theta[\Phi(x)],
\]

Thus, we can write explicitly the FFBRST transformation for the BLG theory as

\[
f \Gamma^{AB}_a = \nabla_a c^{AB}\Theta, \quad f c_{AB} = \frac{1}{2} G^{CD, EF}_{AB} c_{CD} c_{EF}\Theta,
\]

\[
f \bar{\nu}^{AB} = b^{AB}\Theta, \quad f b^{AB} = 0,
\]

\[
f X^{IA} = ic^{AB} X_B^I\Theta, \quad f X^{IA, I} = -i X_B^I c^{AB}\Theta.
\]

The FFBRST transformation is symmetry of the action \( S_{BLG} \) only but not of the generating functional as the Jacobian for path integral measure in the expression of generating functional is not invariant under it. Under FFBRST transformation Jacobian changes as \( \mathcal{D}\Phi^i = J[\Phi(\kappa)]\mathcal{D}\Phi^i(\kappa) \). It has been shown that this nontrivial Jacobian can be replaced within the functional integral as

\[
J[\Phi(\kappa)] \rightarrow e^{iS_1[\Phi(\kappa)]},
\]

where \( S_1[\Phi(\kappa)] \) is some local functional of \( \Phi^i \). The condition for existence of \( S_1 \) is

\[
\int d^3x d^2\theta \left[ \frac{1}{J(\kappa)} \frac{dJ(\kappa)}{d\kappa} - i \frac{dS_1}{d\kappa} \right] = 0.
\]
To calculate the infinitesimal change in Jacobian we use the following expression,

\[
\frac{1}{J(\kappa)} \frac{dJ(\kappa)}{d\kappa} = - \int d^3x d^2\theta \left[ s\Gamma_a(x) \frac{\delta \epsilon[\Phi(x,k)]}{\delta \Gamma_a(x,k)} - sc(x,k) \frac{\delta \epsilon[\Phi(x)]}{\delta c(x,k)} 
- s\tau(x,k) \frac{\delta \epsilon[\Phi(x,k)]}{\delta \tau(x,k)} + sb(x,k) \frac{\delta \epsilon[\Phi(x,k)]}{\delta b(x,k)} \right].
\]  

(24)

5 Relating Different Gauges

In this section, we show explicitly that how FFBRST transformation can be used to analyse the BLG theory in two different gauges. It is possible to take different gauges for the BLG theory. For example, we can take a non-linear gauge in the BLG theory, which is similar to a non-linear gauge in Yang-Mills theories. The sum of the gauge fixing and ghost terms for this non-linear gauge can be written as

\[
\mathcal{L}_{gh} + \mathcal{L}_{gf} = \int d^2 \theta \quad f^{ABMN} \left[ b^{ABD} \Gamma_{aMN} \right. \\
+ \alpha b^{AB} c^{MN} + \frac{1}{2} C^{CD,EF} b_{AB} \right] \\
+ \frac{1}{8} C^{CD,EF} c^{IJ,KL} c^{CD,EF} c^{IJ,KL}.
\]  

(25)

The non-linear BRST transformation are now given by

\[
s\tau^{AB} = b^{AB} - \frac{1}{2} C^{CD,EF} c^{CD,EF} ,
\]

\[
s\Gamma_a^{AB} = \left[ \nabla_a c \right]^{AB} ,
\]

\[
sb^{AB} = - \frac{1}{2} C^{CD,EF} c^{CD,EF} - \frac{1}{8} C^{CD,EF} c^{LM,NP} c^{CD,LM} c^{CD,LP} ,
\]

\[
s\epsilon^{AB} = - \frac{1}{2} C^{CD,EF} c^{CD,EF} ,
\]

\[
s X^I_A = ic^{AB} X^I_B ,
\]

\[
s X^I_A^ \dagger = - i X^I_B^ \dagger c^{AB} .
\]  

(26)

Just like in the linear case here again we can show that these transformations are nilpotent, \( s^2 = 0 \). The non-linear finite BRST transformations are given by

\[
f\tau^{AB} = b^{AB} \Theta - \frac{1}{2} C^{CD,EF} c^{CD,EF} \Theta ,
\]

\[
f\Gamma_a^{AB} = \left[ \nabla_a c \right]^{AB} \Theta ,
\]

\[
f b^{AB} = - \frac{1}{2} C^{CD,EF} c^{CD,EF} \Theta - \frac{1}{8} C^{CD,EF} c^{LM,NP} c^{CD,LM} c^{CD,LP} \Theta ,
\]

\[
f c^{AB} = - \frac{1}{2} C^{CD,EF} c^{CD,EF} \Theta ,
\]
Let the linear and the non-linear gauges be represented by \( G_1^{AB}[\Gamma] \) and \( G_2^{AB}[\Gamma] \) and let \((sG_1)^{AB}\) and \((sG_2)^{AB}\) be the linear BRST and non-linear BRST transformations of these gauge fixing conditions, respectively. The infinitesimal field dependent BRST parameter is now chosen to be

\[
\epsilon[\Phi] = i\gamma \int d^3xd^2\theta \left[ f^{ABCD}c_{AB}(G_{CD1} - G_{CD2}) \right],
\]

where \( \gamma \) is an arbitrary constant parameter. Using expression for the change in Jacobian for this \( \epsilon[\Phi] \) is calculated as

\[
\frac{1}{J} \frac{dJ}{dk} = i\gamma \int d^3xd^2\theta f^{ABCD}[b_{AB}G_{CD1} - b_{AB}G_{CD2} - (sG_{CD1} - sG_{CD2})e_{AB}],
\]

\[
= i\gamma \int d^3xd^2\theta f^{ABCD}[b_{AB}G_{CD1} - b_{AB}G_{CD2} + \bar{c}_{AB}(sG_{CD1} - sG_{CD2})].
\]

Now, an ansatz for \( S_1 \),

\[
S_1 = \int d^3xd^2\theta f^{ABCD}[\xi_1(\kappa)b_{AB}G_{CD1} + \xi_2(\kappa)b_{AB}G_{CD2} + \xi_3(\kappa)\bar{e}_{ABS}GCD1 + \xi_4(\kappa)\bar{e}_{ABS}GCD2],
\]

where arbitrary parameters \( (\xi_i(i = 1, 2, 3, 4)) \) and all fields \( (b^a, \Gamma^a, c^a, \bar{c}^a) \) depend on \( \kappa \). The parameters \( \xi_i(\kappa) \) also satisfy the following initial boundary condition

\[
\xi_i(\kappa = 0) = 0.
\]

Therefore, using Eq. (19) the differentiation of the above equation w. r. to \( \kappa \) gives

\[
\frac{dS_1}{dk} = \int d^3xd^2\theta f^{ABCD}[\xi_1' b_{AB}G_{CD1} + \xi_1 b_{AB}G_{CD1} + \xi_2 b_{AB}G_{CD2}]
\]

\[
+ \xi_2 b_{AB}G_{CD2} + \xi_3' \bar{e}_{ABS}GCD1 - \xi_3 \bar{e}_{ABS}GCD1 + \xi_4 \bar{e}_{ABS}GCD2
\]

\[
+ \xi_4 \bar{e}_{ABS}GCD2 - \xi_4' \bar{e}_{ABS}GCD2,
\]

\[
= \int d^3xd^2\theta f^{ABCD}[\xi_1' b_{AB}G_{CD1} + \xi_2 b_{AB}G_{CD2}]
\]

\[
+ (\xi_1 - \xi_3) b_{AB}G_{CD1} + (\xi_2 - \xi_4) b_{AB}G_{CD2}. \tag{32}
\]

To write the Jacobian \( J \to e^{iS_1} \) the following condition (as mentioned in Eq. (23)) is to be satisfied,

\[
\int d^3xd^2\theta \left[ f^{ABCD}[(\xi_1' - \gamma) b_{AB}G_{CD1} + (\xi_2' + \gamma) b_{AB}G_{CD2} + (\xi_3' - \gamma) \bar{e}_{ABS}GCD1 + (\xi_4' + \gamma) \bar{e}_{ABS}GCD2 + (\xi_1 - \xi_3) b_{AB}G_{CD1} + (\xi_2 - \xi_4) b_{AB}G_{CD2} = 0. \tag{33}
\]
Thus, we get $\xi_1' - \gamma = 0$, $\xi_2' + \gamma = 0$, $\xi_3' - \gamma = 0$, $\xi_4' + \gamma = 0$, $\xi_1 - \xi_3 = 0$, $\xi_2 - \xi_4 = 0$, on equating the coefficients of the above expression. The solutions of above equations satisfying initial condition given in Eq. (31) for $\gamma = 1$ are

$$\xi_1 = \kappa, \xi_2 = -\kappa, \xi_3 = \kappa, \xi_4 = -\kappa.$$  \hspace{1cm} (34)

Now, by adding $S_1(\kappa = 1)$ to the original action having gauge condition $G_{CD2}$ and corresponding ghost term, we get the final action in other gauge $G_{CD1}$ as $S_f = S_{BLG} + S_1$. Thus we see that, $Z$ transforms under FFBRST transformations to

$$Z_f = \int D\Gamma DC\bar{c} Db \ e^{iS_f},$$  \hspace{1cm} (35)

which is nothing but the generating functional for BLG theory in other gauge $G_{CD1}$. This connection is also true for reverse manner i.e. when one starts the theory with gauge condition $G_{CD1}$ and then FFBRST formulation changes the theory in the gauge condition $G_{CD2}$ with appropriate ghost terms.

Thus, the FFBRST formulation can be used to analyse the BLG theory in two different gauges.

6 Conclusion

In this paper we have analysed the BLG theory in $\mathcal{N} = 1$ superspace formalism. As theory had gauge degrees of freedom, we fixed a gauge to quantize it. This gauge fixing condition was incorporated at a quantum level by adding gauge fixing and ghost terms to the original classical Lagrangian. The new effective Lagrangian thus obtained was invariant under a new set of transformations called the BRST transformations. We explicitly wrote the BRST transformations for the BLG theory. These BRST transformations were integrated out to construct the FFBRST transformations. These transformations were constructed by constructing a functional with odd Grassmann parity on the gauge, ghosts, anti-ghosts and auxiliary fields. It did not depend on spacetime explicitly. FFBRST transformations were also found to be the symmetry of the effective action. However, FFBRST transformations did not leave the path integral measure invariant thus were shown to connect the generating functionals of two different effective field theories with suitable choice of the finite field dependent parameter.

It is known that for gauge theories in non-linear gauge the effective Lagrangian is invariant under a larger algebra called Nakanishi-Ojima algebra \cite{17}. It is also known that this algebra is broken by ghost condensation \cite{33, 34}. It will be interesting to analyse the existence of Nakanishi-Ojima algebra and its subsequent breaking in the BLG theory. It may be noted that even for regular Yang-Mills theories the FFBRST for non-linear gauges...
have not been analysed before. It will be interesting to analyse the FF-BRST transformation for both regular Yang-Mills theory and the BLG theory, when the Nakanishi-Ojima algebra is broken by ghost condensation.

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