Topological Massive Gauge Theories in Three Dimensions Based on the Faddeev-Jackiw Formalism

Honglo Lee, Yong-Wan Kim, and Young-Jai Park
Department of Physics,
Sogang University, C.P.O. Box 1142, Seoul 100-611
(Received 29 August 1996)

Abstract

We quantize the (2+1)-dimensional self-dual and Maxwell-Chern-Simons theories by using the Faddeev-Jackiw formulation and compare the results with those of the Dirac formalism.

PACS: 03.70.+k, 11.10.Ef
I. INTRODUCTION

The basic ideas of quantization of a constrained system were first presented by Dirac [1]. By using his method, one can obtain the Dirac brackets, which are the bridges to the commutators in quantum theory. Several years ago, Faddeev and Jackiw (FJ) proposed a method of symplectic quantization of constrained systems for a first-order Lagrangian [2], which was different from the Dirac procedure. In the FJ method, the classification of constraints as first or second class, primary or secondary, is not necessary. All constraints are held to the same standard. Since their work, their quantization method has attracted much attention because it seems to be algebraically much simpler than the Dirac method.

In addition, the study of gauge theories in three-dimensional space-time is very attractive. In odd-dimensional space-time, the topologically non-trivial, gauge-invariant Chern-Simons term gives rise to masses for the gauge fields [3]. It is known that the spin-one theory in 2+1 dimensions may be described by two covariant actions; one is the Maxwell-Chern-Simons (MCS) action, which is constructed with a Maxwell term and a Chern-Simons term [4], while the other is the self-dual (SD) action generating the square root of the Proca equation for a massive vector field [5].

In this paper, we will use both the Dirac and the FJ methods in order to quantize the (2+1)-dimensional gauge theories. We will derive the Dirac brackets and the equivalent equations of motion for the SD model [3,6] in Section II, and for the MCS theory [4,7] in Section III. Section IV presents the conclusion.

II. SD MODEL

1. Dirac Quantization of the SD Model

In this subsection, we first briefly recapitulate the Dirac method with the SD Lagrangian, which is constructed with both the ordinary and the topological mass terms:

$$\mathcal{L} = \frac{1}{2} m^2 B_\mu B^\mu - \frac{1}{2} m \epsilon_{\mu\nu\rho} B^\mu \partial^\nu B^\rho$$  \hspace{1cm} (1)
where $g_{\mu\nu} = diag(1,-1,-1)$ and $\epsilon_{012} = 1$. Denoting the canonical momenta of the vector field as $\Pi_\mu$, we obtain three primary constraints and the canonical Hamiltonian as follows

$$\omega_0 \equiv \Pi_0 \approx 0,$$
$$\omega_i \equiv \Pi_i + \frac{1}{2} m \epsilon_{ij} B^j \approx 0; \quad (i = 1, 2),$$
$$H_c = \int d^2x \left[ -\frac{1}{2} m^2 B^\mu B_\mu + m \epsilon^{ij} B_0 \partial_i B_j \right]. \quad (2)$$

With these primary constraints and the corresponding Lagrange multipliers $\lambda^\mu$, we write the primary Hamiltonian as

$$H_p = H_c + \int d^2x \lambda^\mu \omega_\mu; \quad (\mu = 0, 1, 2). \quad (3)$$

Then, we obtain one more constraint by requiring the time stability of $\omega_0$:

$$\omega_3 \equiv \dot{\omega}_0 = \{\omega_0, H_p\}$$
$$= m^2 B^0 - m \epsilon^{ij} \partial_i B_j \approx 0. \quad (4)$$

The time stabilities of $\omega_i$ and $\omega_3$ give no additional constraints and only play the role of fixing the values of Lagrange multipliers. All four constraints are completely second-class constraints.

According to the Dirac formalism, we can find the $C_{\mu\nu}$-matrix from the Poisson bracket of the constraints

$$C_{\mu\nu} = \{\omega_\mu, \omega_\nu\} = m \begin{pmatrix}
0 & 0 & 0 & -m \\
0 & 0 & -1 & \partial_2^x \\
0 & 1 & 0 & -\partial_1^x \\
m & \partial_2^x & -\partial_1^x & 0
\end{pmatrix} \delta^2(x - y) \quad (5)$$

with the inverse matrix

$$C_{\mu\nu}^{-1} = -\frac{1}{m^2} \begin{pmatrix}
0 & \partial_1^x & \partial_2^x & -1 \\
\partial_1^x & 0 & -m & 0 \\
\partial_2^x & m & 0 & 0 \\
1 & 0 & 0 & 0
\end{pmatrix} \delta^2(x - y). \quad (6)$$
Imposing all the constraints, the reduced Hamiltonian is found to the

\[ H_r = \int d^2 x \left[ \frac{1}{2} (\epsilon^{ij} \partial_i B_j)^2 - \frac{1}{2} m^2 B_i B^i \right] \]

(7)
in which the only physical variables are the \( B_i \). On the other hand, since the Dirac bracket of two variables is defined as

\[ \{A, B\}_D = \{A, B\} - \{A, \omega_\mu\} C_{\mu\nu}^{-1} \{\omega_\nu, B\}, \]

(8)
the non-trivial Dirac brackets of the variables in this model are

\[ \{B_i(x), B_j(y)\}_D = -\frac{1}{m} \epsilon_{ij} \delta^2(x - y). \]

(9)

2. FJ Quantization of the SD Model

Now, we quantize the SD model following the FJ method [2,8,9]. The first-order Lagrangian equivalent to the Lagrangian in Eq.(1) is

\[ \mathcal{L}_{SD} = \frac{m}{2} \epsilon_{ij} \dot{B}^i \dot{B}^j + \mathcal{H}^{(0)}(\xi) \]

(10)
where the zeroth-iterated symplectic potential is

\[ \mathcal{H}^{(0)}(\xi) \equiv m \epsilon_{ij} B^0 \partial_i B^j - \frac{1}{2} m^2 B_u B^u. \]

With the initial set of symplectic variables, \( \xi^{(0)i} = (B^0, B^1, B^2) \), we have, according to the FJ method, the canonical one-form \( a^{(0)}_i = (0, -\frac{m}{2} B^2, \frac{m}{2} B^1) \). These result in the following singular symplectic two-form matrix:

\[ f^{(0)}_{ij}(x, y) = m \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & -1 & 0 \end{pmatrix} \delta^2(x - y). \]

(11)

Note that this matrix has a zero mode, \( \tilde{v}^{(0)}_k(x) = (v_1(x), 0, 0) \), where \( v_1(x) \) is an arbitrary function. From this zero mode, we get the following constraint \( \Omega^{(0)} \):
\[ 0 = \int d^2 x \, \tilde{v}_k^{(0)} \frac{\delta}{\delta \xi^{(0)k}(x)} \int d^2 y \, \mathcal{H}^{(0)}(\xi) \]
\[ = \int d^2 x \, v_1(x) \left[ m\epsilon_{ij} \partial^i B^j - m^2 B^0 \right] \]
\[ \equiv \int d^2 x \, v_1(x) \Omega^{(0)} . \tag{12} \]

In order to provide a consistent description of the system for this constraint, the constrained manifold must be stable under time evolution. In fact, this constraint is stable under time evolution.

According to the FJ method, we can write the first-iterated Lagrangian with a new Lagrange-multiplier as follows:
\[ \mathcal{L}^{(1)} = \frac{m}{2} \epsilon_{ij} B^i \dot{B}^j - \frac{m}{2} B^2 \dot{B}^1 + \Omega^{(0)} \dot{\alpha} - \mathcal{H}^{(1)}(\xi) \tag{13} \]
where the first-iterated symplectic potential is
\[ \mathcal{H}^{(1)}(\xi) = \frac{1}{2} (m B^0)^2 - \frac{1}{2} m^2 B_i \dot{B}^i . \tag{14} \]
Then, the first-iterated set of symplectic variables becomes \( \xi^{(1)i} = (B^0, B^1, B^2, \alpha) \), and the canonical one-form becomes \( a_i^{(1)} = (0, -\frac{m}{2} B^2, \frac{m}{2} B^1, m\epsilon_{ij} \partial^i B^j - m^2 B^0) \). We get the following first-iterated symplectic matrix from the above variables:
\[ f_{ij}^{(1)}(x, y) = m \begin{pmatrix} 0 & 0 & 0 & -m \\ 0 & 0 & 1 & -\partial_x^2 \\ 0 & -1 & 0 & \partial_x^1 \\ m & -\partial_x^2 & \partial_x^1 & 0 \end{pmatrix} \delta^2(x - y). \tag{15} \]
Since this is a non-singular matrix, we finally obtain the desired inverse matrix of the above matrix as
\[ [f_{ij}^{(1)}]^{-1}(x, y) = \frac{1}{m^2} \begin{pmatrix} 0 & -\partial_x^1 & -\partial_x^2 & 1 \\ -\partial_x^1 & 0 & -m & 0 \\ -\partial_x^2 & m & 0 & 0 \\ -1 & 0 & 0 & 0 \end{pmatrix} \delta^2(x - y). \tag{16} \]
Finally, according to the FJ method, the Dirac brackets are acquired directly from the elements of the inverse of the symplectic matrix because

$$\{\xi_i^{(1)}(x), \xi_j^{(1)}(y)\} = [f^{(1)}]_{ij}^{-1}(x, y).$$  \hspace{1cm} (17)

Reading the Dirac brackets from above matrix, we find that

$$\{B_i(x), B_j(y)\}_D = -\frac{1}{m}\epsilon_{ij}\delta^2(x - y),$$  \hspace{1cm} (18)

which is the same as the equation for the Dirac brackets in Eq.(9). In addition, using the constraints and the Dirac brackets, we can easily obtain the self-dual equation of motion for $B_1$, which has only one dynamical degree of freedom,

$$\Box + m^2)B_1 = 0,$$  \hspace{1cm} (19)

because $\Pi_1$ is proportional to $B_2$.

It seems appropriate to comment on the Dirac and the FJ formalisms. Firstly, through the quantization of the SD model, we have shown that the number of constraints is fewer and the structure of these constraints is very simple because we do not need to distinguish between first- or second-class constraints, primary or secondary constraints, etc. Secondly, we have easily obtained the Dirac brackets by reading them directly from the inverse matrix $f^{ij}(x, y)$ of the symplectic two-form matrix. Thirdly, we have shown that the symplectic Hamiltonian at the final stage of iterations exactly gives the reduced physical Hamiltonian, which may be obtained through several steps with the three definitions of the canonical, the total, and the reduced Hamiltonians in the usual Dirac formulation for constrained systems.

The above three merits have been recently analyzed in several papers, on the subjects of the nonrelativistic point particle, three-dimensional topologically massive electrodynamics, the nonlinear sigma model, two-dimensional induced gravity [8], constrained systems [9], etc. These works show how efficient the symplectic formalism is, and confirm that the symplectic quantization method is a simpler alternative to the Dirac’s formalism in the sense that the brackets are obtained more easily and are exactly same as the Dirac brackets. As a result, we can replace the obtained brackets with the quantum commutators as $\{ , \}_D \rightarrow i[ , ]$. 

6
III. MCS THEORY

1. Dirac Quantization of the MCS Theory

In this subsection, in order to compare it with the FJ formalism, we sketch the Dirac quantization procedure with the MCS theory, which is constructed with the Maxwell and the topological mass terms:

\[ \mathcal{L}_{MCS} = -\frac{1}{2} F_{\mu} F_{\mu} + \frac{1}{2} m F_{\mu} A_{\mu} \]  

(20)

where \( F_{\mu} \equiv \frac{1}{2} \epsilon_{\mu\nu\rho} F^{\nu\rho} = \epsilon_{\mu\nu\rho} \partial^\nu A^\rho \). Denoting the canonical momenta of the vector field as \( \Pi_\alpha \), we obtain one primary constraint and the canonical Hamiltonian as follows:

\[ \omega_0 = \Pi_0 \equiv 0, \]  

(21)

\[ H_c = \int d^2 x [ -\frac{m}{2} \epsilon_{ij} \Pi_i A_j + \frac{1}{2} (\Pi_i)^2 + \frac{m^2}{8} (A^i)^2 + \frac{1}{2} (\epsilon_{ij} \partial^j A^i)^2 + \partial_i \partial^j A^0 - \frac{m}{2} \epsilon_{ij} (\partial^j A^i) A_0 ] . \]  

(22)

With the primary constraint and the corresponding Lagrange multiplier \( u \), we write the primary Hamiltonian as

\[ H_p = H_c + \int d^2 x u \omega_0. \]  

(23)

Requiring time stability of the primary constraint, we get one more constraint:

\[ \omega_1 \equiv \dot{\omega}_0 = \partial^i \Pi_i + \frac{m}{2} \epsilon_{ij} \partial^j A_j. \]  

(24)

Note that the time stability of \( \omega_1 \) gives no additional constraints and only plays the role of fixing the value of the Lagrange multiplier. These two constraints are first class, which gives rise to gauge invariance. Therefore, we should introduce a gauge-fixing function to find the true physical variables correctly. Choosing the Coulomb gauge condition \( \omega_2 = \partial_i A^i \), we obtain one more constraint:

\[ \omega_3 \equiv \dot{\omega}_2 = m \epsilon_{ij} \partial^j A^i + \partial_i \partial^i A^0. \]  

(25)
Now, all four constraints are second-class.

We find the $C_{\mu\nu}$-matrix from the Poisson bracket of the constraints:

$$C_{\mu\nu} = \nabla^2 \begin{pmatrix} 0 & 0 & 0 & -1 \\ 0 & 1 & 0 & 0 \\ -1 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 \end{pmatrix} \delta^2(x - y)$$

(26)

with its inverse

$$C^{-1}_{\mu\nu} = \frac{1}{\nabla^2} \begin{pmatrix} 0 & 0 & 0 & 1 \\ 0 & 0 & -1 & 0 \\ 0 & 1 & 0 & 0 \\ -1 & 0 & 0 & 0 \end{pmatrix} \delta^2(x - y).$$

(27)

Imposing all the constraints on Eq. (23), the reduced Hamiltonian is found to be

$$H_r = \int d^2x \left[ -\frac{m}{2} \epsilon_{ij} \Pi_i A^j + \frac{1}{2} (\Pi_i)^2 + \frac{m^2}{8} (A^i)^2 + \frac{1}{2} (\epsilon_{ij} \partial^i A^j)^2 \right].$$

(28)

Through a similar procedure as in the previous section, we obtain the following Dirac brackets:

$$\{\Pi_i(x), \Pi_j(y)\}_D = -\frac{m}{2} \epsilon_{ij} \delta^2(x - y),$$

$$\{A^i(x), \Pi_j(y)\}_D = \frac{\epsilon^{ik} \epsilon_{ij} \partial^k \partial^l \delta^2(x - y)}{\nabla^2}.$$

(29)

These Dirac brackets will be compared with the symplectic brackets in the next subsection.

2. FJ Quantization of the MCS Theory

Now, we quantize the MCS theory following the FJ method. The first-order Lagrangian is

$$\mathcal{L}_{MCS} = \Pi_i \dot{A}^i - H^{(0)}(\xi)$$

(30)

where the zeroth-iterated symplectic potential is
\[ H^{(0)}(\xi) = -\frac{m}{2} \epsilon_{ij} \Pi_i A^j + \frac{1}{2} (\Pi_i)^2 + \frac{m^2}{8} (A^i)^2 + \frac{1}{2} (\epsilon_{ij} \partial^j A^i)^2 + \Pi_i \partial^i A^0 - \frac{m}{2} \epsilon_{ij} (\partial^i A^j) A_0. \] (31)

With the initial set of symplectic variables, \( \xi^{(0)i} = (A^0, A^1, A^2, \Pi_1, \Pi_2) \), we have, according to the FJ method, the canonical one-form \( a^{(0)}_i = (0, \Pi_1, \Pi_2, 0, 0) \). These result in the following singular symplectic two-form matrix:

\[
  f^{(0)}_{ij}(x, y) = \begin{pmatrix}
    0 & 0 & 0 & 0 & 0 \\
    0 & 0 & -1 & 0 & 0 \\
    0 & 0 & 0 & -1 & \delta^2(x - y) \\
    0 & 1 & 0 & 0 & 0 \\
    0 & 0 & 1 & 0 & 0
  \end{pmatrix}
\] (32)

This matrix has a zero mode \( \tilde{v}_k^{(0)}(x) = (v_1(x), 0, 0, 0, 0) \), where \( v_1(x) \) is an arbitrary function. Using this zero mode, we get the following constraint:

\[
  0 = \int d^2 x \frac{\delta}{\delta \xi^{(0)k}(x)} \int d^2 y \ H^{(0)}(\xi) = -\int d^2 x \ v_1(x) [\partial^i \Pi_i + \frac{m}{2} \epsilon_{ij} \partial^j A^i] \equiv -\int d^2 x \ v_1(x) \Omega^{(0)}. \] (33)

We can write the first-iterated Lagrangian with a new Lagrange-multiplier as

\[
  \mathcal{L}^{(1)}_{MCS} = \Pi_i \dot{A}^i + \Omega^{(0)} \dot{\alpha} - H^{(1)}(\xi) \] (34)

where the first-iterated symplectic potential is

\[
  H^{(1)}(\xi) = H^{(0)}(\xi) \mid_{\Omega^{(0)}=0} = -\frac{m}{2} \epsilon_{ij} \Pi_i A^j + \frac{1}{2} (\Pi_i)^2 + \frac{m^2}{8} (A^i)^2 + \frac{1}{2} (\epsilon_{ij} \partial^j A^i)^2. \] (35)

Then, the first-iterated set of symplectic variables becomes \( \xi^{(1)i} = (A^1, A^2, \Pi_1, \Pi_2, \alpha) \), and the canonical one-form becomes \( a^{(1)}_i = (\Pi_1, \Pi_2, 0, 0, \partial_i \Pi_i + \frac{m}{2} \epsilon_{ij} \partial^j A^i) \). From these variables, we find the following first-iterated symplectic matrix:
This matrix is also singular.

Although we use the zero mode, $\tilde{v}^{(1)}_k = (\partial_1 v_5, \partial_2 v_5, \frac{m}{2} \partial_2 v_5, -\frac{m}{2} \partial_1 v_5, v_5)$, which gives non-dynamical relations of the system in the FJ method, we can’t obtain the constraint any more. Since the Lagrangian one-form is invariant under the transformation rule of the symplectic variable, $\delta \xi^{(1)i} = \tilde{v}^{(1)}_k \eta$, we should introduce a gauge-fixing function. Using the Coulomb gauge condition $\Omega^{(1)} = \partial^i A_i$, we can extend the system as follows

$$L^{(2)}_{MCS} = \Pi_i \dot{A}_i + \Omega^{(0)} \dot{\alpha} + \Omega^{(1)} \dot{\beta} - H^{(2)}(\xi)$$

where

$$H^{(2)}(\xi) = H^{(1)}(\xi) |_{\Omega^{(1)} = 0}.$$ 

The symplectic variables and the canonical one-form of the second-iterated Lagrangian are

$$\xi^{(2)i} = (A^1, A^2, \Pi_1, \Pi_2, \alpha, \beta),$$

$$\sigma^{(2)i} = (\Pi_1, \Pi_2, 0, 0, \partial_i \Pi^i + \frac{m}{2} \epsilon_{ij} \partial^j A^i, \partial_i A^i).$$

Then, the symplectic two-form matrix is

$$f^{(2)}_{ij}(x, y) = \begin{pmatrix}
0 & 0 & -1 & 0 & -\frac{m}{2} \partial_x^2 - \partial_x^1 \\
0 & 0 & 0 & -1 & \frac{m}{2} \partial_x^1 - \partial_x^2 \\
0 & 0 & 0 & 0 & \partial_x^1 \\
1 & 0 & 0 & 0 & \partial_x^2 \\
-\frac{m}{2} \partial_x^2 & \frac{m}{2} \partial_x^1 & \partial_x^1 & \partial^2 & 0 & 0 \\
-\partial_x^1 & -\partial_x^2 & 0 & 0 & 0 & 0
\end{pmatrix} \delta^2(x - y).$$
Since this matrix is non-singular, we finally obtain inverse as

\[
[f^{(2)}_{ij}]^{-1}(x, y) = \frac{1}{\nabla^2} \begin{pmatrix}
0 & 0 & \partial_x^2 \partial_y^2 & -\partial_x^1 \partial_y^2 & 0 & \partial_y^1 \\
0 & 0 & -\partial_x^1 \partial_y^2 & \partial_x^1 \partial_y^1 & 0 & \partial_y^2 \\
\partial_x^2 \partial_y^2 & \partial_x^1 \partial_y^2 & 0 & -\frac{m}{2} & \partial_x^1 & \frac{m}{2} \partial_y^2 \\
\partial_x^1 \partial_y^2 & \partial_x^1 \partial_y^1 & \frac{m}{2} & 0 & -\partial_x^2 & -\frac{m}{2} \partial_y^1 \\
0 & 0 & -\partial_x^1 & -\partial_x^2 & 0 & 1 \\
\partial_x^1 & \partial_x^2 & \frac{m}{2} \partial_y^2 & -\frac{m}{2} \partial_y^1 & -1 & 0
\end{pmatrix} \delta^2(x - y). \tag{40}
\]

Then, we can directly read the Dirac brackets for the true physical fields from the above matrix, and they are the same as those in Eq. (29). In addition, we know that the physical degree of freedom in the configuration space is only one. The equation of motion of this degree of freedom, which is really the dual field \(F^0\) contained in the Lagrangian in Eq. (20), is obtained by using the Dirac brackets in Eq. (29) and is found to be

\[(\Box + m^2)F^0 = 0. \tag{41}\]

Therefore, the field \(F^0\) is effectively equivalent to the field \(B_1\), with the same mass appearing in both Eqs. (19) and (41).

**IV. CONCLUSION**

In conclusion we have studied the MCS theory and the SD gauge theory in (2+1)-dimensions using the Dirac and the FJ formulations. We have found that both the Dirac and the FJ formulations result in the same Dirac brackets. Especially, we ascertain that for these cases the FJ formulation also is algebraically a much simpler method, which gives the desired Dirac brackets readily without the classification of constraints, than that of Dirac’s just as several other interesting models \[\text{seealso}\]. We have also shown that both the MCS theory and the SD gauge theory have only one degree of freedom in the configuration space and have effectively the same equations of motion. From this fact, we have found through FJ quantization that both theories are equivalent to each other at the level of the equation of
motion. This result coincides with that of Faddeev and Jackiw obtained by using the Master equation.

ACKNOWLEDGMENTS

The present study was supported by the Basic Science Research Institute Program, Ministry of Education, Project No. BSRI-96-2414.
REFERENCES

[1] P. A. M. Dirac, *Lectures on Quantum Mechanics* (Belfer Graduate School, Yeshiva University Press, New York, 1964).

[2] L. Faddeev and R. Jackiw, Phys. Rev. Lett. **60**, 1692 (1988).

[3] W. Siegel, Nucl. Phys. **B156**, 135 (1979); R. Jackiw and S. Templeton, Phys. Rev. **D23**, 2291 (1981).

[4] S. Deser, R. Jackiw, and S. Templeton, Ann. Phys. **140**, 372 (1982).

[5] P. K. Townsend, K. Pilch, and P. van Nienwenhuizen, Phys. Lett. **B136**, 38 (1984).

[6] Y. W. Kim, Y. J. Park, K. Y. Kim, and Y. Kim, Phys. Rev. **D51**, 2943 (1995).

[7] S. Deser and R. Jackiw, Phys. Lett. **B139**, 371 (1984).

[8] M. M. Horta Barreira and C. Wotzasek, Phys. Rev. **D45**, 1410 (1992); J. Barcelos-Neto and C. Wotzasek, Mod. Phys. Lett. **A7**, 1737 (1992); Int. J. Mod. Phys. **A7**, 4981 (1992); C. Wotzasek and C. Neves, J. Math. Phys. **34**, 1807 (1993); C. Han, Phys. Rev. **D47**, 5521 (1993).

[9] D. S. Kulshreshtha and H. J. W. Müller-Kirsten, Phys. Rev. **D45**, R393 (1992); N. Banerjee, D. Chatterjee, and S. Ghosh, Phys. Rev. **D46**, 5590 (1992); Y-W. Kim, Y-J. Park, K. Y. Kim, Y. Kim, and C-H. Kim, J. Korean Phys. Soc. **26**, 243 (1993); J. W. Jun and C. Jue, Phys. Rev. **D50**, 2939 (1994); S-J. Yoon, Y-W. Kim, S-K. Kim, Y-J. Park, K.Y. Kim, and Y. Kim, J. Korean Phys. Soc. **27**, 270 (1994); Y-W. Kim, Y-J. Park, K.Y. Kim, and Y. Kim, J. Korean Phys. Soc. **27**, 610 (1994); Y-W. Kim, Y-J. Park, and Y. Kim, J. Korean Phys. Soc. **28**, 773 (1995); E-B. Park, Y-W. Kim, Y-J. Park, and Y. Kim, Mod. Phys. Lett. **A10**, 1119 (1995).