<table>
<thead>
<tr>
<th>Title</th>
<th>BPS operators from the Wilson loop in the 3-dimensional supersymmetric Chern-Simons theory</th>
</tr>
</thead>
<tbody>
<tr>
<td>Author(s)</td>
<td>Fujita, Mitsutoshi</td>
</tr>
<tr>
<td>Citation</td>
<td>Physical Review D - Particles, Fields, Gravitation and Cosmology (2009), 80(8)</td>
</tr>
<tr>
<td>Issue Date</td>
<td>2009-10</td>
</tr>
<tr>
<td>URL</td>
<td><a href="http://hdl.handle.net/2433/89517">http://hdl.handle.net/2433/89517</a></td>
</tr>
<tr>
<td>Rights</td>
<td>© 2009 The American Physical Society.</td>
</tr>
<tr>
<td>Type</td>
<td>Journal Article</td>
</tr>
<tr>
<td>Textversion</td>
<td>publisher</td>
</tr>
</tbody>
</table>

Kyoto University
BPS operators from the Wilson loop in the 3-dimensional supersymmetric Chern-Simons theory

Mitsutoshi Fujita*

Department of Physics, Kyoto University, Kyoto 606-8502, Japan
(Received 20 July 2009; published 2 October 2009)

We consider the small deformation of the pointlike Wilson loop in the 3-dimensional \( \mathcal{N} = 6 \) superconformal Chern-Simons theory. By Taylor expansion of the pointlike Wilson loop in powers of the loop variables, we obtain the BPS operators that correspond to the excited string states of the dual IIA string theory on the \( pp \) wave background. The BPS conditions of the Wilson loop constrain both the loop variables and the forms of the operators obtained in the Taylor expansion.

DOI: 10.1103/PhysRevD.80.086001 PACS numbers: 11.25.Tq, 11.30.Pb

I. INTRODUCTION

The dual gravity interpretation of the supersymmetric Wilson loop in the \( D = 4 \) \( \mathcal{N} = 4 \) super Yang-Mills (SYM) theory has been important in the context of the AdS/CFT correspondence [1]. The supersymmetric Wilson loop contains the 6 scalar fields \( \Phi^i \) and the expectation value of the Wilson loop is protected from the UV divergence [2]. The expectation value of the circular Wilson loop was obtained by analyzing the string minimal surface in the anti–de Sitter space [3]. Furthermore, the supersymmetric Wilson loop in the Berenstein-Maldacena-Nastase (BMN) sector can be described by the dual IIB string theory on the \( pp \) wave background. The BMN correspondence [4] is the duality between the infinite strings of operators in the dual IIB string theory on the \( pp \) wave background. In Ref. [5], it was conjectured that the 1/2-Bogomol’nyi-Prasad-Sommerfield (BPS) pointlike Wilson loop \( W(C_0) \) for \( C_0 \) shrinking to a spacetime point is mapped to the vacuum state of the dual IIB string theory on the \( pp \) wave background. The functional derivatives of the Wilson loop are mapped to the excited string states. We have summarized the results obtained in Ref. [5] in Table I.

In this paper, motivated by the work in Ref. [5], we consider the dual IIA string theory description of the supersymmetric Wilson loop in the recently proposed \( D = 3 \) \( \mathcal{N} = 6 \) Chern-Simons-matter theory [Aharony-Bergman-Jafferis-Maldacena (ABJM) theory] [6,7]. The ABJM theory is the low-energy effective theory of the \( N \) M2-branes at the singularity of the orbifold \( \mathbb{C}^4/\mathbb{Z}_k \).

This theory can be analyzed by using the dual IIA string theory on the \( \text{AdS}_4 \times \mathbb{C}P^3 \) spacetime and on its Penrose limit [8–12] in the parameter regime

\[
\sqrt{\lambda} \gg 1, \quad e^{2\phi} \sim \frac{N^{1/2}}{k^{3/2}} \ll 1, \quad (1.1)
\]

where \( k \) is the Chern-Simons level, \( N \) is the rank of the gauge group, \( \lambda = N/k \) is \( \chi \) Hooft coupling, and \( \phi \) is the dilaton. The first one in (1.1) implies that, in the dual type IIA string theory, the radius of curvature is much larger than 1 in the string unit, and the second one implies that we take the small string coupling limit, to suppress the quantum corrections.

The supersymmetric Wilson loop in the ABJM theory was proposed in the literature [13–15]. The Wilson loop contains a product of the bifundamental scalars on the exponent. It was shown that the straight line and circular Wilson loops preserve 1/6 of the ABJM supersymmetry.

The main purpose of our paper is to study the dual IIA string theory description of the pointlike Wilson loop that has enhanced 1/3 supersymmetry. We show that by deforming the pointlike Wilson loop, we can obtain the BPS operators that correspond to the excited string states of the dual IIA string theory on the \( pp \) wave background. The BPS conditions of the Wilson loop, (3.10) and (3.11), give the constraint on both the loop variables and the forms of the BPS operators. By following the conjecture in Ref. [5], we give maps from the functional derivatives of the Wilson loop to the dual IIA string excited states.

The Penrose limit of the dual gravity theory is given by the following limit:

\[
N, J \to \infty \quad \text{with} \quad \lambda' = \frac{\lambda}{J^2} \text{ fixed,} \quad (1.2)
\]

where \( J \) is the charge of the infinite strings of operators under the \( U(1) \) subgroup of \( SU(4) \) R symmetry. By determining the function \( h(\lambda) \) that appears in the dispersion relation, the gauge/gravity correspondence has been proved up to the curvature corrections to the \( pp \) wave background [17].

The content of this paper is as follows: In Sec. II, we consider the pointlike Wilson loop in the ABJM theory and compare it with the vacuum state of the dual IIA superstring theory. In Sec. III, we obtain the BPS conditions for the Wilson loop. We show that the pointlike Wilson loop is 1/3 BPS and the supersymmetry generator preserved by the pointlike Wilson loop is the same as that preserved by the infinite chain dual to the IIA string vacuum state. In

---

*mfujita@gauge.scphys.kyoto-u.ac.jp

1In the special case of the gauge group \( SU(2) \times SU(2) \), the ABJM theory has \( SO(8) \) enhanced global R symmetry.

2See [16] on the deformation of the Wilson loop operator in the \( \mathcal{N} = 4 \) SYM as well as in the YM theory.
TABLE I. The relation between the functional derivatives and the IIB string oscillation modes $\alpha^{(a)}_n$ and $\alpha^{(a)}_n$. In the table, $\mu = 0, 1, 2, 3, a = 1, 2, 3, 4$, and $Z = \Phi^5 + i\Phi^6$.

<table>
<thead>
<tr>
<th>SYM side</th>
<th>Dual IIB string side</th>
</tr>
</thead>
<tbody>
<tr>
<td>Functional derivatives</td>
<td>Corresponding operators</td>
</tr>
<tr>
<td>$\cdot \cdot \cdot$</td>
<td>$W(C_0)$</td>
</tr>
<tr>
<td>$\int d\mathcal{S} e^{2\pi i n/\mu} \frac{\delta}{\delta \mathcal{S}(s)}$</td>
<td>$D_{\mu}Z$</td>
</tr>
<tr>
<td>$\int d\mathcal{S} e^{2\pi i n/\mu} \frac{\delta}{\delta \mathcal{S}(\Phi^6)}$</td>
<td>$\Phi^6$</td>
</tr>
</tbody>
</table>

Sec. IV, we solve the BPS equations satisfied by the loop variables and expand the Wilson loop in powers of the independent loop variables. Thus, we obtain the maps from the functional derivatives of the Wilson loop to the dual IIA string excited states.

II. THE SUPERSYMMETRIC WILSON LOOP IN THE ABJM THEORY

The ABJM theory is the 3-dimensional $\mathcal{N} = 6$ supersymmetric Chern-Simons theory with the gauge group $U(N) \times U(N)$. The fields in the ABJM theory are the $U(N) \times U(N)$ gauge fields $A_m$ and $\tilde{A}_m$, the bifundamental bosonic fields $Y^I$ [14], and the bifundamental spinor fields $\psi_{IJ}$, where $I = \{1, 2, \ldots, 4\}$ is the index of $SU_R(4)$ $R$ symmetry and $\alpha$ ($\alpha = 1, 2$) is the (2 + 1)-dimensional spinor index.

According to Ref. [13], the supersymmetric Wilson loop in the ABJM theory is given by

$$W[C] = \text{Tr} \left[ P \exp \left( \int_C ds \left( \hat{x}^m(s)A_m + M^J_f(s)Y^I_fY^J_f \right) \right) \right]$$

(2.1)

where $\hat{x}^m(s)$ describes the path $C$ on $R^2$ and the function $M_f^J(s)$, determined by the supersymmetry (SUSY), will be the coordinates of the transverse space $\mathbb{C}^4/\mathbb{Z}_k$. We assume that $M_f^J(s)$ is a $4 \times 4$ real matrix.5

In Ref. [14], they also obtain the Wilson loop with the gauge field $A_m$. We do not consider the Wilson loop with $\tilde{A}_m$ here; our Wilson loop breaks the parity symmetry of the ABJM theory (see [11,18]).

We can give $M_f^J(s)$ a $U_R(1)$ charge. $U_R(1)$ is a subgroup of $SU_R(4)$ $R$ symmetry. In the Wilson loop (2.1), the $U_R(1)$ symmetry that rotates $A_1$ and $B_1$ by $\alpha = \exp(\frac{\pi i}{2})$ also operates on $M_f^J(s)$ as follows:

$$V^{-1}M = \begin{pmatrix} \alpha & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & \alpha & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} m_{11} & m_{12} & m_{13} & m_{14} \\ m_{21} & m_{22} & m_{23} & m_{24} \\ m_{31} & m_{32} & m_{33} & m_{34} \\ m_{41} & m_{42} & m_{43} & m_{44} \end{pmatrix} \begin{pmatrix} \alpha & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & \alpha & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} \tilde{m}_{11} & \tilde{m}_{12} & \tilde{m}_{13} & \tilde{m}_{14} \\ \tilde{m}_{21} & \tilde{m}_{22} & \tilde{m}_{23} & \tilde{m}_{24} \\ \tilde{m}_{31} & \tilde{m}_{32} & \tilde{m}_{33} & \tilde{m}_{34} \\ \tilde{m}_{41} & \tilde{m}_{42} & \tilde{m}_{43} & \tilde{m}_{44} \end{pmatrix}$$

In Sec. IV, we use this $U_R(1)$ symmetry to classify the loop variables.

We consider the pointlike Wilson loop whose path $C_0$ shrinks to the point $x^m = x_0^m$ ($x^m = 0$). We set $M^J_f$ in a nilpotent matrix,

$$\begin{pmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}$$

(2.2)

By expanding the exponential part of the Wilson loop, we obtain the infinite sum of the local operator as follows:

$$W[C_0] = \text{Tr} \left[ \exp(itA_1B_1(x_0)) \right] = \sum_{j=0}^{\infty} \frac{(it)^j}{j!} \text{Tr}[(A_1B_1)^j](x_0),$$

(2.3)

where $t$ describes the periodicity of the loop; we identify $s = 0$ with $s = t$. In higher order of $J$, Eq. (2.2) includes the infinite strings of the operator $A_1B_1$ that correspond to the vacuum state of the dual IIA string theory. In Sec. IV, by deforming (2.3), we also obtain the BPS operators that correspond to the dual IIA string excited states.

In next section, we analyze the SUSY preserved by the Wilson loop (2.1) and (2.3).

III. SUPERSYMMETRY

In Sec. III A, we shortly review the SUSY transformation of the supersymmetric Wilson loop in the ABJM theory. In Sec. III B, we derive the BPS conditions satisfied by the supersymmetric Wilson loop. In Sec. III C, we show that the pointlike Wilson loop (2.3) is 1/3 BPS.

A. Supersymmetry transformation of the Wilson loop

The space of the 6 Majorana spinors, which are also the $\mathcal{N} = 6$ SUSY generator. The $\mathcal{N} = 6$ SUSY transformations are given by

$$\delta Y^I = i\omega^{IJ}_a \psi_J,$$

(3.3)

$$\delta Y^I = i\omega^{IJ}_a \psi_J,$$

(3.4)

$$\delta A_m = -Y^I \psi_J \gamma^m \omega_{IJ} + \omega_{IJ} \gamma^m \psi_I Y^J_f,$$

(3.5)

$$\delta A_m = -Y^I \psi_J \gamma^m \omega_{IJ} + \omega_{IJ} \gamma^m \psi_I Y^J_f.$$
where the convention of the spinors is the same as given in Refs. [7,19].

We consider the SUSY transformation of the Wilson loop as follows:

\[ e^{i(\omega_{ij} Q^{ij} + \omega^{ij} \Pi^{ij})} W[C] e^{-i(\omega_{ij} Q^{ij} + \omega^{ij} \Pi^{ij})} = W[C] + \delta W[C], \]

(3.7)

where \( Q^{ij} \) is the SUSY generator. The Wilson loop preserves a part of the SUSY in the ABJM theory when \( \delta W[C] = 0 \) for an arbitrary \( s \).

By using the condition \( \delta W[C] = 0 \), we can show that the SUSY generator \( \omega_{AB} \) preserved by the Wilson loop satisfies the following equations as given in Ref. [13]:

\[ \omega_{AB}^{\alpha} \gamma_{\alpha \beta} \dot{x}^\beta - i M^{(K)}_{K} \omega_{KA, \beta} = 0, \]  

(3.8)

\[ \omega_{AB, \alpha} \gamma_{\alpha \beta} \dot{x}^\beta - i M^{(A)}_I \omega^{IB} = 0. \]  

(3.9)

Note that the complex conjugate of (3.8) gives (3.9) when \( M^{(I)}_I \) is a Hermitian matrix.

B. BPS conditions

When we contract (3.8) and (3.9) by \( \epsilon_{\alpha \beta} \) and the charge conjugation matrix \( \tilde{C} \) (see Appendix A) or when we multiply (3.8) by \( \gamma_{\alpha} \dot{x}^\alpha - i M^T \) from the right, we obtain the following BPS conditions:

\[ 4 \dot{x}^2(s) + M^{(K)}_{K}(s) M^{(I)}_I(s) = 0, \]  

(3.10)

\[ \det(\dot{x}^2(s) + M^T(s)M(s)) = \det(\dot{x}^2(s) + M(s)M(s)) = 0. \]  

(3.11)

(3.10) and (3.11) are the necessary condition to preserve a part of SUSY: The straight line and (2.3) satisfy (3.10). Equation (3.10) is similar to the BPS conditions \( \dot{x}^2 + \dot{y}^2 = 0 \) satisfied by the Wilson loop in the \( d = 4 \), \( \mathcal{N} = 4 \) SYM.

C. Supersymmetry preserved by the pointlike Wilson loop

We can show easily that the pointlike Wilson loop (2.3) preserves 1/3 of the SUSY in the ABJM theory. For \( \dot{x}_m = 0 \), Eqs. (3.8) and (3.9) are given by

\[ M^{(K)}_{K} \omega_{KA, \beta} = 0, \]  

(3.12)

\[ M^{(I)}_I \omega^{IB} = 0. \]  

(3.13)

By solving these equations, we obtain \( \omega_{13} = \omega_{23} = \omega_{24} = \omega_{34} = 0 \). The pointlike Wilson loop preserves the SUSY \( \omega_{12} \) and \( \omega_{14} \) which are not constrained by (3.12) and (3.13). Note that the Wilson loop over the straight path is 1/6 BPS. We guess that there is an enhancement of the SUSY when we shrink the loop to a point.

We want to explain why the Wilson loop in the ABJM theory is 1/3 BPS or 1/6 BPS instead of 1/2 BPS. In the dual IIA string theory side, the Wilson loop is described by the fundamental string on \( AdS_5 \times CP^3 \) spacetimes. From the supersymmetry analysis [14] of the Killing spinors, it has been known that the fundamental string dual to the straight Wilson loop is not localized at \( CP^3 \) but smeared along \( CP^3 \): The smeared string preserves less SUSY. So, we guess that a similar phenomenon happens for our pointlike Wilson loop.

IV. BPS OPERATORS FROM THE WILSON LOOP

In this section, we show that the BPS operators arise in the double series expansion of the pointlike Wilson loop operator (2.3) in powers of the loop variables \( \delta x^m(s) \) and \( \delta M^I_I(s) = m_{IJ}(s) \) (see also [5]).

First, we consider the Wilson loop fluctuated near the point \( x^m(s) = x_0^m \). We parametrize \( M^I_I(s) \) by

\[ M^I_I(s) = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix} + \begin{pmatrix} m_{11} & m_{12} & 0 & m_{14} \\ m_{21} & m_{22} & 0 & m_{24} \\ m_{31} & m_{32} & m_{33} & m_{34} \\ m_{41} & m_{42} & m_{43} & m_{44} \end{pmatrix}, \]  

(4.1)

where we fix the gauge freedom of the parameter \( s \) by imposing \( M_1 \), and for convenience, omit the label \( s \) in \( m_{IJ}(s) \). Since the loop coordinates \( \delta x^m(s) \) and \( m_{IJ}(s) \) should be periodic about \( s \), they can be rewritten as follows:

\[ \delta x^m(s) = \sum_{n=-\infty}^{\infty} \delta x_0^m e^{2\pi ins/t}, \]  

(4.2)

\[ m_{IJ}(s) = \sum_{n=-\infty}^{\infty} m_{nIJ} e^{2\pi ins/t} \quad \text{for} \ (I, J) \neq (3, 1). \]

In the Fourier space, the loop variables have zero modes.

The loop variables \( \delta x^m(s) \) and \( \delta M^I_I(s) \) are not independent but are related by the BPS condition (3.10) and (3.11). By substituting (4.1) into the matrix \( X_{ab} = (\delta x^2 + M^2)_{ab} \) and the BPS condition (3.10) and (3.11), we obtain the following matrix elements and the following equation:

\[ \frac{1}{2} \left( X_{00} - X_{11} - X_{22} - X_{33} + X_{44} \right) = 0. \]
\[ X_{11} = \delta \dot{\chi}^2 + m_{11}^2 + m_{12}m_{21} + m_{13} + m_{14}m_{41}, \quad X_{12} = m_{11}m_{12} + m_{12}m_{22} + m_{32} + m_{14}m_{42}, \]
\[ X_{13} = m_{11} + m_{12}m_{23} + m_{33} + m_{14}m_{43}, \quad X_{14} = m_{11}m_{14} + m_{12}m_{24} + m_{34} + m_{14}m_{44}, \]
\[ X_{21} = m_{21}m_{11} + m_{22}m_{21} + m_{23}m_{31} + m_{24}m_{41}, \quad X_{22} = \delta \dot{\chi}^2 + m_{12}m_{21} + m_{22}^2 + m_{23}m_{32} + m_{24}m_{42}, \]
\[ X_{23} = m_{21} + m_{22}m_{23} + m_{23}m_{33} + m_{24}m_{43}, \quad X_{24} = m_{21}m_{14} + m_{22}m_{24} + m_{23}m_{34} + m_{24}m_{44}, \]
\[ X_{31} = m_{31}m_{11} + m_{32}m_{21} + m_{33}m_{31} + m_{34}m_{41}, \quad X_{32} = m_{31}m_{12} + m_{32}m_{22} + m_{33}m_{32} + m_{34}m_{42}, \]
\[ X_{33} = \delta \dot{\chi}^2 + m_{31} + m_{32}m_{33} + m_{33}^2 + m_{34}m_{43}, \quad X_{34} = m_{31}m_{14} + m_{32}m_{24} + m_{33}m_{34} + m_{34}m_{44}, \]
\[ X_{41} = m_{41}m_{11} + m_{42}m_{21} + m_{43}m_{31} + m_{44}m_{41}, \quad X_{42} = m_{41}m_{12} + m_{42}m_{22} + m_{43}m_{32} + m_{44}m_{42}, \]
\[ X_{43} = m_{41} + m_{42}m_{23} + m_{43}m_{33} + m_{44}m_{43}, \quad X_{44} = \delta \dot{\chi}^2 + m_{41}m_{44} + m_{24}m_{42} + m_{34}m_{43} + m_{44}^2, \]

\[ \det X = X_{11}X_{22}X_{33}X_{44} - X_{11}X_{22}X_{34}X_{43} + X_{11}X_{23}X_{34}X_{42} - X_{11}X_{23}X_{32}X_{44} - X_{11}X_{24}X_{32}X_{43} - X_{11}X_{24}X_{34}X_{42} - X_{11}X_{24}X_{32}X_{43} = 0. \]

A special solution of (4.4) is given by\(^8\)
\[ m_{22} = m_{44} = m_{24} = m_{42} = 0, \]
\[ m_{32} = -m_{11}m_{12}, \quad m_{34} = -m_{11}m_{14}, \quad m_{41} = -m_{43}m_{33}, \]
\[ m_{21} = -m_{23}m_{33}, \quad m_{11} = -m_{33}, \]
\[ \delta \dot{\chi}^2 = 0, \]
\[ \Rightarrow X_{12} = X_{14} = X_{43} = X_{23} = X_{42} = X_{24} = X_{22} = X_{44} = 0. \]

In the Fourier space, the relations (4.5), (4.6), (4.7), (4.8), (4.9), and (4.10) are rewritten as follows:
\[ m_{32(n)} = -m_{11(n)}m_{12(-n)}, \quad m_{34(n)} = -m_{11(n)}m_{14(-n)}, \]
\[ m_{41(n)} = -m_{43(n)}m_{33(-n)}, \]
\[ m_{21(n)} = -m_{23(n)}m_{33(-n)}, \quad m_{11(n)} = -m_{33(n)}, \]
\[ \delta \chi_n^0 \delta \chi_n^0 = \delta \chi_n^1 \delta \chi_n^1 + \delta \chi_n^2 \delta \chi_n^2 \quad \text{for} \ n \neq 0, \]
\[ m_{31(n)} = m_{11(n)}m_{11(-n)}. \]

Note that, while the zero modes $\delta \chi_n^0$ are the independent parameters, $\delta \chi_n^0$ is not independent.

Recall that the loop variable has the $U_p(1)$ charges as given in Table II (see footnote 4). The loop variable also has the dimension $\Delta$ since the dimension of the Wilson loop becomes zero. Since the charge $Q(-\Delta - J)$ of the loop variables is well-defined for (4.6), (4.7), (4.8), (4.9), (4.10), we expand the Wilson loop fluctuated at $x_n^0$ up to one in powers of the charge $Q$ of the loop variable and up to one impurity as follows:
\[ W[C] = W[C_0] + \int_0^t ds \sum_{Q(m_{ij})=-1} m_{ij}(s) \frac{\delta W[C]}{\delta m_{ij}(s)} \bigg|_{C=C_0} + \int_0^t ds \chi_n^m(s) \frac{\delta W[C]}{\delta \chi_n^m(s)} \bigg|_{C=C_0} + \cdots, \]

where we used (4.6), (4.7), and (4.8). Here, $\sum_{Q(m_{ij})=-n}$ means that we sum the terms in which the charge of $m_{ij}$ is $n$; for $n = 1/2$, we sum the terms containing $m_{12}$, $m_{14}$, $m_{23}$, and $m_{43}$.

We introduce a new functional derivative as follows:
\[ \frac{\delta}{\delta m'} = \frac{\delta}{\delta m_{11}} - \frac{\delta}{\delta m_{33}}. \]

By substituting the Fourier transformation of the loop variables (4.2) into the resulting formula, we obtain
\[ W[C] = W[C_0] + \sum_{Q(m_{ij})=-1/2} \sum_n m_{ij}(n) \int_0^t ds \frac{\delta W[C]}{\delta m_{ij}(s)} \bigg|_{C=C_0} \times e^{2i\pi n s/t} + \sum_n m_{11(n)} \int_0^t ds \frac{\delta W[C]}{\delta m'(s)} \bigg|_{C=C_0} e^{2i\pi n s/t} + \cdots. \]

The functional derivative of (2.3) contained in (4.17) generates the impurity operator (in the sense of the spin chain)
as follows [20,21]:

\[
\frac{\delta W(C)}{\delta x^m(s)} \bigg|_{C\to C_0} = i \text{Tr}[(F_{mn}(x(s))\delta s + (D_m Y^l Y^l_j) \\
\quad \times (\delta x(s))M^l_j(s)w^s_i(C))\bigg|_{C\to C_0} = i \sum_{j=0}^{\infty} \frac{(it)^j}{j!} \text{Tr}[(D_mAIB_1(x_0)(AIB_1(x_0))^j],
\]

(4.18)

\[
\frac{\delta W(C)}{\delta M^L(s)} \bigg|_{C\to C_0} = i \sum_{j=0}^{\infty} \frac{(it)^j}{j!} \text{Tr}[Y^k Y^l_i(x_0)(AIB_1(x_0))^j],
\]

(4.19)

where \( w^s_i(C) \) describes the Wilson line along the path \( C \) from \( s = s' \) to \( s = s'' \) and the covariant derivative \( D_m \) including \( A_m \) operates on \( Y^l Y^l_j \) so that \( Y^l Y^l_j \) is an adjoint field. Note that (4.18) and (4.19) do not depend on the loop parameter \( s \).

By using (4.18) and (4.19), we find that the right side of (4.17) becomes the zero mode contribution as follows:

\[
\sum_{j=0}^{\infty} \frac{(it)^j}{j!} \left[ \text{Tr}[(AIB_1)^j](x_0) \\
+ it \sum_{Q(m_0)=-1/2} m_{ij0} \text{Tr}[Y^i Y^l_j(AIB_1)^j](x_0) \\
+ itm_{i10} \text{Tr}[(Y^i Y^l_1 - Y^3 Y^l_3)(AIB_1)^j](x_0) \\
+ it\delta x^m_0 \text{Tr}[(D_mAIB_1(AIB_1)^j](x_0)] \right].
\]

(4.20)

In higher order of \( J \), Eq. (4.20) contains the following infinite strings of the operator \( AIB_1 \), the bifundamental scalars \( Y^l Y^l_j \), and the covariant derivative \( D_m \):

\[
\text{Tr}[(AIB_1)^j](x_0), \quad \text{Tr}[Y^i Y^l_j(AIB_1)^j](x_0), \quad \text{Tr}[(D_mAIB_1(AIB_1)^j](x_0),
\]

(4.21)

\[
\text{Tr}[(Y^i Y^l_1 - Y^3 Y^l_3)(AIB_1)^j](x_0),
\]

(4.22)

where \( (i, j) = (1, 2, 1), (1, 4, 2), (2, 3, 4) \) and we separated the operators coupled to the independent loop variables. These BPS operators protected by the supersymmetry correspond to the vacuum state and the excited states of the dual IIA string theory on the \( pp \) wave background (see [8,9]).

By following the correspondence between (4.21) and the dual IIA string excited states, and by following the conjecture in Ref. [5], we map the functional derivatives of the Wilson loop to the dual IIA string excited states in Table III.

Before ending this section, we want to comment on the relation between our Taylor expansion of the pointlike Wilson loop and strong coupling expansion: The Wilson loop should be expanded in terms of \( 1/\sqrt{s} \) [22] and have a convergent region of the expansion. The operators obtained in the Taylor expansion should be normalized by using their 2-point function. However, the appearance of strong coupling expansion is not clear in both our pointlike Wilson loop and that of \( \mathcal{N} = 4 \) SYM [5]. On the other hand, our Taylor expansion of the Wilson loop (4.20) is similar to the mode expansion of the wave function in the string field theory (SFT) as was pointed out in the work of Ref. [5]. Since the SFT is independent of the string coupling, (4.20) may be independent of the ’t Hooft coupling.

V. DISCUSSIONS

In this paper, motivated by the paper [5], we discussed the dual IIA string description of a slightly deformed pointlike Wilson loop. By expanding the pointlike Wilson loop in powers of the loop variables, we obtained the BPS operators that correspond to the excited string states of the IIA string theory on the \( pp \) wave background. Our new result was the impurity operator \( Y^l Y^l_j - Y^3 Y^3_j \) that was not determined from the analysis in the gravity side [8]. By following the conjecture in Ref. [5], we gave the maps from the functional derivatives of the Wilson loop to the dual IIA string excited states in Table III.

The BPS conditions of the Wilson loop (3.10) and (3.11) constrained both the loop variables and the forms of the BPS operators, though we did not prove the uniqueness of solution of the BPS conditions. In Appendix C, we also obtained the BMN operators [23] in higher order terms of the Taylor expansion (C1). It will also be important to compute the anomalous dimension [24] of the BMN operators (C15) and (C16).

Since our Wilson loop is pointlike, it is important to construct the straight line BPS Wilson loop that connects the pointlike 1/3-BPS Wilson loop to the 1/6-BPS Wilson loop in the ABJM theory. We leave it to future work.

We also want to comment on the Wilson loop in the ABJM theory from the viewpoint of 11-dimensional theory.

<table>
<thead>
<tr>
<th>( m_{13} )</th>
<th>( m_{12}, m_{14}, m_{23}, m_{43} )</th>
<th>( \delta x^m )</th>
<th>( m_{1i}, m_{2i}, m_{3i}, m_{4i} )</th>
<th>( m_{21}, m_{41}, m_{32}, m_{34} )</th>
<th>( m_{31} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>1/2</td>
<td>0</td>
<td>-1/2</td>
<td>-1</td>
<td>1</td>
</tr>
</tbody>
</table>

---

9As was pointed out in Ref. [9], we need to include the fermions to complete the BPS multiplet that the impurity operators belong to, though the gauge/gravity duality could be proven by using only the bosonic string action.
TABLE III. The relation between the functional derivatives and the IIA supergravity modes (the zero modes of the string oscillation operator). In the table, $m = 0, 1, 2$, $(i, j) = (1, 2), (1, 4), (2, 3), (4, 3)$, and $k = 5, 6, 7, 8$.

<table>
<thead>
<tr>
<th>ABJM side</th>
<th>Dual IIA superstring side</th>
</tr>
</thead>
<tbody>
<tr>
<td>Dual string</td>
<td>IIA Corresponding operators</td>
</tr>
<tr>
<td>$\ldots$</td>
<td>$W(C_0)$</td>
</tr>
<tr>
<td>$\int ds e^{2\phi_{mns}^{i}} \delta_{mns}^{j}$</td>
<td>$D_{n}a_{s}B_{j}$</td>
</tr>
<tr>
<td>$\int ds e^{2\phi_{mns}^{i}} \delta_{mns}^{j}$</td>
<td>$Y^{i}Y^{j} - Y^{j}Y^{i}$</td>
</tr>
<tr>
<td>$\int ds e^{2\phi_{mns}^{i}} \delta_{mns}^{j}$</td>
<td>$Y^{i}Y^{j}$</td>
</tr>
</tbody>
</table>

(M theory). It has been known that in the dual IIA supergravity side, the Wilson loop is described by the fundamental string on the AdS$_4 \times$ CP$^3$ spacetime. Then, the dual fundamental string should be described by the M2-brane wrapped on the M-circle: The 3-form $C_{(3)}$ in the 11-dimensional supergravity reduces to the Neveu-Schwarz–Neveu-Schwarz field $B_{\mu\nu}$ coupling to the IIA string. Thus, it is also interesting to analyze the dynamics of the Wilson loop by using the loop equation [25]. As the dynamics of the Wilson loop describe the dynamics of the dual string, the loop equation will also describe the dynamics of the M2-brane wrapped on the M-circle. When we analyze the loop equation, however, we need to extend the analysis to include supersymmetry (see also [21]).

ACKNOWLEDGMENTS

Especially, we thank M. Fukuma, H. Hata, H. Kawai, and K. Yoshida for the discussions and helpful comments. We thank K. Katayama for the discussions. We thank T. Nishioka and T. Takayanagi for the helpful comments on this manuscript.

APPENDIX A: CHARGE CONJUGATION MATRIX

In Appendix A, we introduce the charge conjugation matrix of both SO(6) and SO(4) and construct the reducible representation of SO(6) gamma matrices by using the reducible representation of SO(4) gamma matrices. The charge conjugation matrix constructed by SO(6) gamma matrices satisfies the following relations:

$$C\gamma^{i}C^{-1} = -\gamma^{iT},$$

$$C^{T} = \eta C,$$

$$\gamma^{i_1i_2\ldots i_n}C^{T} = \eta \gamma^{i_1i_2\ldots i_n}C$$

where (A1) is the definition of the charge conjugation matrix and $\eta$ is a constant determined later. (A2) shows that the matrix defined on the right-hand side of (A2) transforms as the symmetric tensor or the antisymmetric tensor in terms of the SO(6) spinor index. Since the product representation in Table IV must be consistent with the decomposition $8_+ \otimes 8_+ \rightarrow 36_2 \oplus 28_4$, $\eta$ must be $+1$.

$C$ defined in (A1) is the metric of the spinor index as follows:

$$C = C_{\alpha\beta} = C^{\gamma}, \quad C^{-1} = (C^{-1})^{\alpha\beta}$$

$$C_{\alpha\beta}(\gamma^{i})^{\beta \gamma} = (C\gamma^{i})_{\alpha \gamma},$$

$$C\gamma^{i} = (C\gamma^{i})^{T}, \quad C_{\alpha\beta}\gamma^{a\gamma}g^{\beta \delta} = C\gamma^{\delta}.$$  

where $g$ is the spinor representation of SO(6).

We introduce the SO(4) charge conjugation matrix $\hat{C}$ that satisfies the following relations:

$$\hat{C}\rho^{k}\hat{C}^{-1} = -(\rho^{k})^{T},$$

$$\hat{C}^{T} = -\hat{C}.$$  

We decompose the SO(6) Dirac gamma matrices and the charge conjugation matrix as follows:

$$\gamma^{i} = \begin{pmatrix} 0 & \rho^{i} \\ \rho^{k} & 0 \end{pmatrix}, \quad \gamma^{5} = \begin{pmatrix} 0 & \rho^{5} \\ \rho^{5} & 0 \end{pmatrix}, \quad \gamma^{6} = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}$$

$$C = \begin{pmatrix} 0 & \hat{C} \\ -\hat{C} & 0 \end{pmatrix}.$$  

where $\rho^{k}$ is the SO(4) Dirac gamma matrix. Here the matrices defined in (A7) and (A8) satisfy the Clifford algebra and the definition of the charge conjugation matrix.

APPENDIX B: CLEBSCH-GORDAN DECOMPOSITION OF THE 6 MAJORANA SPINORS

We show the Clebsch-Gordan decomposition of the 6 SO(1,2) Majorana spinors $\epsilon$, which transform as the vector representation of spin(6) $\sim$ SO(6), and show the equality (3.2). We know the following equivalence:

$$\text{spin (6)} \sim SU(4),$$

$$6_\epsilon \sim 6_\alpha,$$

$$4_\epsilon \sim 4_\alpha.$$

We introduce the new matrix $C_2$ and 4$_\epsilon$ spinor index I and J as follows:

\footnote{These equations are satisfied for any SO(N).}

\footnote{We can show the second equation in (A4) by expanding the exponential of $g$ in powers of the generator of the Lorentz algebra.}
TABLE IV. The product representation of the BPS operators from the Wilson loop in the 3-…

<table>
<thead>
<tr>
<th>n</th>
<th></th>
<th>0</th>
<th>1</th>
<th>2</th>
<th>3</th>
<th>4</th>
<th>5</th>
<th>6</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\delta C_n$</td>
<td>$(-1)^{(1/2)n(n+1)}$</td>
<td>1</td>
<td>6</td>
<td>15</td>
<td>20</td>
<td>15</td>
<td>6</td>
<td>1</td>
</tr>
</tbody>
</table>

By using (A4), we can show that

$$W_{i} = \sum_{I<J} (\hat{C}_2 \hat{\gamma}^i)_{IJ} v^{IJ},$$

where

$$\gamma^i = \begin{pmatrix} 0 & \hat{\gamma}^i_J \\ \hat{\gamma}^i_I & 0 \end{pmatrix}.$$ 

By using (A4), we can show that $\hat{C}_2 \hat{\gamma}^i$ transforms as an antisymmetric tensor in terms of the index $I, J$.

We describe the Clebsch-Gordan decomposition of $\hat{C}_2 \hat{\gamma}^i$ as follows:

$$v_{IJ} \rightarrow v' = \sum_{I<J} (\hat{C}_2 \hat{\gamma}^i)_{IJ} v^{IJ},$$

We choose the basis of the $SO(4)$ gamma matrices as follows:

$$\rho^1 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \rho^2 = \begin{pmatrix} 0 & \sigma^1 \\ \sigma^1 & 0 \end{pmatrix},$$

and

$$\rho^3 = \begin{pmatrix} 0 & \sigma^3 \\ \sigma^3 & 0 \end{pmatrix}, \quad \rho^4 = \begin{pmatrix} 0 & \sigma^2 \\ \sigma^2 & 0 \end{pmatrix}.$$ 

After a $SO(6)$ permutation of the $SO(6)$ gamma matrices $\gamma^i$ including a change of the signs, we obtain the supersymmetry generator $\omega_{IJ}$, which satisfies (3.1) and (3.2) (see also [19]).

It is convenient to introduce the following equations:

$$(C \gamma^i)^* = -\gamma^i C^*,$$

$$(C \gamma^7)^* = \gamma^7 C^*.$$ 

By using the $SO(6)$ gamma matrices $C$ and $C_2$, we can show (3.2) as follows:

$$\epsilon \epsilon(C_2 \gamma^i) \epsilon(C_2^\dagger \gamma^i) = \epsilon \epsilon(C \gamma^7 \gamma^i C^* \gamma^i)$$

$$= \epsilon \epsilon(C \gamma^i \gamma^7 C^*) = \epsilon \epsilon(C^*)$$

where $\gamma^7 = \gamma^1 \gamma^2 \gamma^3 \gamma^4$.

**APPENDIX C: HIGHER ORDER TERMS IN THE EXPANSION OF THE POINTLIKE WILSON LOOP**

In this section, we compute higher order terms in the expansion (4.15). We expand the pointlike Wilson loop up to 2 in powers of the charge $Q$ as follows:

$$\hat{C} = \rho^2 \rho^3, \quad \hat{C}_2 = \rho^4 \rho^1.$$ 

(9)

$$\hat{C} = \rho^2 \rho^3, \quad \hat{C}_2 = \rho^4 \rho^1.$$ 

(9)
where \( \sum_{Q(m_j \neq m_k) = n} \) means that we sum the terms in which the charge of \( m_j m_k \) is \( n \). Recall that the loop variables satisfy the relation (4.11), (4.12), (4.13), and (4.14) as follows:

\[
m_{32(n)} = -m_{11(n)}m_{12(-n)}, \quad m_{34(n)} = -m_{11(n)}m_{14(-n)},
\]

\[
m_{41(n)} = -m_{43(n)}m_{33(-n)}, \quad (C2)
\]

\[
m_{21(n)} = -m_{23(n)}m_{33(-n)}, \quad m_{11(n)} = -m_{33(n)}. \quad (C3)
\]

Because of (C4), the spacetime coordinates are twisting for \( n \neq 0 \). By substituting the Fourier transformation of the loop variables (4.2) into (C1) and by using (C2)–(C5), we obtain

\[
W[C] = W[0] + \sum_{Q(m_i) = 1/2} \sum_n \sum m_{ij(n)} \int_0^t ds \left[ \frac{\delta W[C]}{\delta m_{ij}(s)} \right] e^{2\pi i n \alpha} + \sum_n \sum_{m_{ij(n)}} m_{ij(n)} \int_0^t ds_1 \left[ \frac{\delta^2 W[C]}{\delta m_{ij}(s_1) \delta m_{ij}(s_2)} \right] e^{2\pi i n \alpha} / n \neq 0 \)

\[
\delta x_n^0 \delta x_{-n} = \delta x_n^1 \delta x_{-n} + \delta x_n^2 \delta x_{-n} \quad \text{for } n \neq 0, \quad (C4)
\]

\[
m_{31(n)} = m_{11(n)}m_{11(-n)}. \quad (C5)
\]

where the loop variables \( \delta x_n^m \) for \( n \neq 0 \) are not independent. Note that, in the context of the spin chain, the impurity \( D_m \) coupling with \( \delta x_n^m \) mixes with the fermions: We guess that the Wilson loop containing the fermions will be needed to explain the full excited string spectrum in the IIA string theory. The two functional derivatives of the pointlike Wilson loop are given by

\[
\frac{\delta^2 W(C)}{\delta x_n^m(s_2) \delta x_n^m(s_1)} \bigg|_{C-C_0} = \delta \left[ \frac{\delta W[C]}{\delta x_n^m(s_2)} \right] i \text{Tr} \left( F_m(x(s_1)) \delta x_n^m(s_1) \right) + D_m Y^l Y^\dag_l (x(s_1)) M_l^l(s_1) w_s^{l+1}(C) \bigg|_{C-C_0}
\]

\[
= i \text{Tr} \left[ D_m A_l B_l(x_0) w_s^{l+1}(C_0) \right] + i \text{Tr} \left[ F_m w_s^{l+1}(C_0) \right] - \text{Tr} \left[ D_m A_l B_l(x_0) w_s^{l+1}(C_0) D_m A_l B_l(x_0) w_s^{l+1}(C_0) \right]. \quad (C7)
\]

\[
\frac{\delta^2 W(C)}{\delta M_{K}^l(s_2) \delta x_n^m(s_1)} \bigg|_{C-C_0} = \delta \left[ \frac{\delta W[C]}{\delta M_{K}^l(s_2)} \right] i \text{Tr} \left( F_m(x(s_1)) \delta x_n^m(s_1) \right) + D_m Y^l Y^\dag_l (x(s_1)) M_l^l(s_1) w_s^{l+1}(C) \bigg|_{C-C_0}
\]

\[
= i \text{Tr} \left[ D_m A_l B_l(x_0) w_s^{l+1}(C_0) \right] - \text{Tr} \left[ D_m A_l B_l(x_0) w_s^{l+1}(C_0) D_m A_l B_l(x_0) w_s^{l+1}(C_0) \right]. \quad (C8)
\]
The derivation of the second line in (C10) is given in the appendix of [5]. By substituting (C7)–(C9) into (C6) and by transforming the parameter \((s_1, s_2)\) into \((s_1/t, (s_2 - s_1)/t)\) for the double integral about \(s_1\) and \(s_2\), we obtain the local operator expression of the remaining terms as follows: The \(Q = 1\) terms in which a product of the loop variables has the charge \(Q = 1\) become

\[
\sum_{Q(m_n, n_m) = 1} \sum_{\mathbf{i}J} (it)^{\mathbf{i}J + 1} \sum_{\mathbf{i}J} \left[ \sum_{k=0}^j \text{Tr} [(Y^1 Y^1_j - Y^3 Y^3_j) (A_1 B_1)^k Y^1_j Y^1_0 (A_1 B_1)^{j-k}] (x_0) \cdot F_2(n_1, k, J) \right] (x_0) F_2(n_1, k, J).
\]

The \(Q = 3/2\) terms become

\[
+ \sum_{\mathbf{i}J} \sum_{\mathbf{i}J} (it)^{\mathbf{i}J + 1} \sum_{\mathbf{i}J} \left[ \sum_{k=0}^j \text{Tr} [(Y^1 Y^1_j - Y^3 Y^3_j) (A_1 B_1)^k Y^1_j Y^1_0 (A_1 B_1)^{j-k}] (x_0) \cdot F_2(n_1, k, J) \right] (x_0) F_2(n_1, k, J)
- \frac{1}{(J + 1)!} \text{Tr} [Y^3 Y^3_j] (A_2 B_1)^{j+1}] (x_0) + \sum_{\mathbf{i}J} \sum_{\mathbf{i}J} (it)^{\mathbf{i}J + 1} \sum_{\mathbf{i}J} \left[ \sum_{k=0}^j \text{Tr} [(Y^1 Y^1_j - Y^3 Y^3_j) (A_1 B_1)^k Y^1_j Y^1_0 (A_1 B_1)^{j-k}] (x_0) \cdot F_2(n_1, k, J) \right] (x_0) F_2(n_1, k, J)
+ \frac{1}{(J + 1)!} \text{Tr} [Y^3 Y^3_j] (A_2 B_1)^{j+1}] (x_0) + \cdots.
\]

where we have not written the terms dependent on \(\delta x^m\) since they are not independent. We obtain the \(Q = 2\) terms up to two impurity as follows:

\[
+ \frac{1}{2} \sum_{\mathbf{i}J} \sum_{\mathbf{i}J} (it)^{\mathbf{i}J + 1} \sum_{\mathbf{i}J} \left[ \sum_{k=0}^j \text{Tr} [(Y^1 Y^1_j - Y^3 Y^3_j) (A_1 B_1)^k Y^1_j Y^1_0 (A_1 B_1)^{j-k}] (x_0) \cdot F_2(n_1, k, J) \right] (x_0) F_2(n_1, k, J)
+ \frac{1}{(J + 1)!} \text{Tr} [Y^3 Y^3_j] (A_2 B_1)^{j+1}] (x_0) + \cdots.
\]

where we have not written the terms dependent on \(\delta x^m\) for the same reason. Note that \(Q = 2\) terms also contain the 3 impurity terms.

In the large \(J\) limit, (C11)–(C13) contain the following BMN operators:
\[ \sum_{k=0}^{j} \text{Tr}[Y^{i} Y_{i}^{j}(A_{1}B_{1})^{k} Y^{j} Y_{n}^{j}(A_{1}B_{1})^{j-k}] e^{2\pi ink/J}, \] (C14)

\[ \sum_{k=0}^{j} \text{Tr}[(Y^{i} Y_{i}^{j} - Y^{3} Y_{3}^{j})(A_{1}B_{1})^{k} Y^{j} Y_{4}^{j}(A_{1}B_{1})^{j-k}] e^{2\pi ink/J} - \text{Tr}[Y^{3} Y_{3}^{j}(A_{1}B_{1})^{j+1}], \] (C15)

\[ \sum_{k=0}^{j} \text{Tr}[(Y^{i} Y_{i}^{j} - Y^{3} Y_{3}^{j})(A_{1}B_{1})^{k} Y^{j} Y_{4}^{j}(A_{1}B_{1})^{j-k}] e^{2\pi ink/J} + \text{Tr}[Y^{3} Y_{3}^{j}(A_{1}B_{1})^{j+1}], \] (C16)

where \( (i, j) \) and \( (l, m) \) are equal to \( (1, 2), (1, 4), (2, 3), \) and \( (4, 3) \).