Holographic estimate of oblique corrections for technicolor

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(Received 24 February 2006; published 19 July 2006)

We study the oblique corrections to the electroweak interaction in the holographic model of technicolor theories. The oblique S parameter is expressed in terms of a solution to the equations of motion for the anti de Sitter bulk gauge fields. By analyzing the solution, we establish a rigorous proof that the S parameter is positive and is reduced by walking. We also present the precise numerical values for the S parameter of various technicolor models by solving the equations numerically.

DOI: 10.1103/PhysRevD.74.015011 PACS numbers: 12.60.Nz, 11.10.Kk, 11.25.Tq, 12.15.Lk

The development of the anti de Sitter/conformal field theory (AdS/CFT) correspondence [1] has attracted a lot of interest in recent years, since it may shed some light on strongly coupled gauge theories, whose understanding has been otherwise quite limited. Inspired by recent success [2,3] in understanding the infrared (IR) dynamics of quantum chromodynamics (QCD) with its holographic dual, we apply the AdS/CFT correspondence to estimate the oblique corrections of technicolor theories.

The precision electroweak data have shown that the standard model of electroweak interaction is extremely stable radiatively, which is often expressed in terms of vanishingly small Peskin-Takeuchi S, T, U parameters [4,5]. Nonzero values of those parameters indicate new physics beyond the standard model. Any model for the new physics is highly constrained by the measured values. Since the oblique correction of technicolor is dominated by the nonperturbative dynamics in the IR region, its precise estimate has been a major hurdle for technicolor theories to be candidates for physics beyond the standard model. Recently, however, a new class of technicolor models (techni-orientifold) [6,7], which uses a higher-dimensional representation for techniquarks and the correspondence with $\mathcal{N} = 1$ super Yang-Mills theory at large N [8], has been proposed. The models need only a small number of technifermions to be in the conformal window, free from the flavor-changing neutral-current problem, and thus have naturally small oblique corrections [7,9,10].

In this article we attempt to calculate precisely the S parameter for technicolor theories, using the AdS/CFT correspondence. (Similar attempts were made in 5-dimensional Higgsless models [11].) We also show rigorously that the S parameter is strictly positive for the holographic dual models of dynamical electroweak symmetry breaking and the walking behavior of the technicolor dynamics reduces the S parameter substantially.

According to AdS/CFT correspondence, to every operators in CFT there correspond AdS bulk fields. The bulk fields that satisfy the equations of motion are the sources for the CFT operators, when evaluated at the ultraviolet (UV) boundary of the AdS space, and their action at the UV boundary is nothing but the generating functional for the connected Green functions of those operators in CFT.

For a near conformal theory whose IR scale is generated far below the UV scale, we may take a slice of the AdS metric as

$$ds^2 = \frac{1}{z^2}(-dz^2 + \eta^{\mu\nu}dx_\mu dx_\nu), \qquad \epsilon \leq z \leq z_m, \quad (1)$$

where $z = \epsilon$ ($z = z_m$) is the UV (IR) boundary, and $\eta^{\mu\nu}$ is the four-dimensional Minkowski metric. For the boundary conformal field theory we choose a strongly coupled technicolor theory with the $SU(N)_{TC}$ gauge group and N_{TF} massless techniquarks, q^α ($\alpha = 1, \ldots, N_{TF}$), which may be fundamental under the technicolor gauge group as in the walking technicolor [12] or second-rank tensor as in the techni-orientifold [6,7]. Since the technicolor model we choose for the boundary field theory is not exactly conformal, the AdS/CFT correspondence does not hold in a strict sense. We expect, however, the holographic dual description of technicolor to work reasonably well, as the conformal symmetry is mildly broken, especially for the technicolor theories in the conformal window. The boundary operators we are interested in are the scalar, $\bar{q}_\alpha q^\beta$, and the (axial) vector current, $J^{a\mu}_{V(A)} = \bar{q}\gamma^\mu(\gamma_5)t^a q$, of techniquarks. [$t^a$'s are the $SU(N_{TF})$ generators, normalized as $\mathrm{Tr} t^a t^b = 1/2\delta^{ab}$.] The corresponding bulk action is then given as

$$S = \int d^5x\sqrt{g}\,\mathrm{Tr}\left[|DX|^2 - m_5^2|X|^2 - \frac{1}{2g_5^2}(F_L^2 + F_R^2)\right] \quad (2)$$

where m_5 is the mass of the bulk scalar field X and the covariant derivative $D_\mu X = \partial_\mu X - iA_{L\mu}X + iXA_{R\mu}$. The bulk mass is determined by the relation [13,14] $\Delta(\Delta - 4) = m_5^2$, where Δ is the dimension of the corresponding boundary operator $\bar{q}_\alpha q^\beta$. F_L and F_R are the field strength tensors of the $SU(N_{TF})_L \times SU(N_{TF})_R$ bulk gauge fields A_L and A_R, respectively. The values of vector and axial gauge

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fields, defined as $V = (A_L + A_R)/\sqrt{2}$ and $A = (A_L - A_R)/\sqrt{2}$, at the UV boundary couple to $J^a_{V\mu}$ and $J^a_{A\mu}$, respectively.

The classical solution of the bulk field X for $\Delta > 2$ is given as $X_0(z) = c_1 z^{4-\Delta} + c_2 z^{\Delta}$, where the constants c_1 and c_2 are determined by the boundary conditions. The chiral condensate of techniquarks is (formally) defined as

$$\langle \bar{q}_L q_R \rangle = i \frac{\delta}{\delta S} e^{iW[S]} \Big|_{S=0} \equiv i \frac{\delta}{\delta S} \langle e^{-i \int_x (\bar{q}_L S q_R + \mathrm{H.c.})} \rangle_{S=0}, \tag{3}$$

where S is the source for the chiral condensate and the generating functional is given by the AdS/CFT correspondence

$$W[S] = \int_x \left(-\frac{1}{4z^3} \partial_z X_0 X_0^\dagger + \mathrm{H.c.} \right)_{z=\epsilon} \tag{4}$$

with the identification $S(x) = 2X_0(x, z)/z^{4-\Delta}|_{z=\epsilon}$. Since the source for the chiral condensate is the (bare) mass, we find the UV boundary condition for X_0,

$$\frac{2}{z^{4-\Delta}} X_0|_{z=\epsilon} = M \quad \text{or} \quad c_1 = \frac{1}{2} M. \tag{5}$$

Then, $c_2 = 4\langle \bar{q} q \rangle / \Delta$ in the chiral limit ($c_1 \to 0$).

The oblique S parameter is defined in terms of the two-point function of the techniquark currents,

$$i \int_x e^{iq\cdot x} \langle J_\mu(x) J_\nu(0) \rangle = -\left(g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2} \right) \Pi(-q^2) \tag{6}$$

as

$$S = 4\pi \frac{d}{dq^2} [\Pi_V(-q^2) - \Pi_A(-q^2)]_{q^2=0}, \tag{7}$$

where Π_V (Π_A) is the (axial) vector current correlator. Since the source for the current is given by the bulk gauge fields, the generating functional for the current correlation functions becomes

$$W[V, A] = -\frac{1}{2g_5^2} \int_x \left(\frac{1}{z} V^a_\mu \partial_z V^{\mu a} + \frac{1}{z} A^a_\mu \partial_z A^{\mu a} \right)_{z=\epsilon}, \tag{8}$$

using the equations of motion. The UV boundary is identified as the extended technicolor (ETC) scale, $\epsilon = 1/\Lambda_{\mathrm{ETC}}$. The IR boundary is given by the technicolor scale, $z_m = 1/\Lambda_{\mathrm{TC}}$, at which all techniquarks get mass and decouple.

The gauge fields in Eq. (8) satisfy the bulk equations of motion in unitary gauge,

$$\begin{aligned} \left[\left(\partial^2 - z\partial_z \frac{1}{z} \partial_z \right) \eta_{\mu\nu} - \partial_\mu \partial_\nu \right] V^\nu = 0, \\ \left[\left(\partial^2 - z\partial_z \frac{1}{z} \partial_z + \frac{g_5^2 X_0^2}{z^2} \right) \eta_{\mu\nu} - \partial_\mu \partial_\nu \right] A^\nu = 0, \end{aligned} \tag{9}$$

where $\partial^2 = \eta^{\mu\nu} \partial_\mu \partial_\nu$ and X_0 is a solution to the equation of motion for the bulk scalar field X, together with suitable boundary conditions that satisfy

$$\sqrt{g}\, \mathrm{Tr}(\delta A_{L\mu} F_L^{\mu z} + \delta A_{R\mu} F_R^{\mu z})|_\epsilon^{z_m} = 0. \tag{10}$$

For simplicity we choose $\delta A_{L\mu} = 0 = \delta A_{R\mu}$ at $z = \epsilon$ and $F_L^{\mu z} = 0 = F_R^{\mu z}$ at $z = z_m$. Furthermore, we will work with the gauge $A_{Lz} = 0 = A_{Rz}$ to have $\partial_z A_{L\mu} = 0 = \partial_z A_{R\mu}$ at $z = z_m$. We have neglected the nonlinear terms for the gauge fields in Eq. (9), since we are interested only in the two-point functions of the gauge fields, and this corresponds to keeping the leading term in the large N_{TC} expansion.

From Eqs. (8) and (9) it is simple to read off the two-point correlation functions of currents in momentum space. We write the 4D Fourier transform of the vector and axial gauge fields as $V^\mu(q, z) = (g^{\mu\nu} - q^\mu q^\nu/q^2) V(q, z)$, $A^\mu(q, z) = (g^{\mu\nu} - q^\mu q^\nu/q^2) A(q, z) + q^\mu q^\nu/q^2 A(0, z)$ with $V(q, \epsilon) = A(q, \epsilon) = 1$ and $\partial_z V(q, z_m) = \partial_z A(q, z_m) = 0$, where $V(q, z)$ and $A(q, z)$ satisfy

$$\begin{aligned} \left[z\partial_z \left(\frac{1}{z} \partial_z \right) + q^2 \right] V(q, z) = 0, \\ \left[z\partial_z \left(\frac{1}{z} \partial_z \right) + q^2 - \frac{g_5^2 X_0^2}{z^2} \right] A(q, z) = 0. \end{aligned} \tag{11}$$

Then we have

$$\begin{aligned} \Pi_V(-q^2) = \frac{\partial_z V(q, z)}{z g_5^2} |_{z=\epsilon}, \\ \Pi_A(-q^2) = \frac{\partial_z A(q, z)}{z g_5^2} |_{z=\epsilon}. \end{aligned} \tag{12}$$

The solution for $V(q, z)$ is given as, with $|q| = \sqrt{|q^2|}$,

$$V(q, z) = a_1 |q| z Y_1(|q| z) + a_2 |q| z J_1(|q| z), \tag{13}$$

where J_1 and Y_1 are the first order Bessel and Neumann functions, respectively. The constants a_1 and a_2 are to be fixed by the boundary conditions $V(q, \epsilon) = 1$ and $\partial_z V(q, z_m) = 0$. From this we have

$$\Pi_V(-q^2) = \frac{|q|}{g_5^2 \epsilon} \frac{J_0(|q| z_m) Y_0(|q| \epsilon) - Y_0(|q| z_m) J_0(|q| \epsilon)}{J_0(|q| z_m) Y_1(|q| \epsilon) - Y_0(|q| z_m) J_1(|q| \epsilon)}, \tag{14}$$

which becomes $q^2 \ln(z_m/\epsilon)/g_5^2$ as $q^2 \to 0$. Comparing it also with the perturbative calculation at large momentum, $-q^2 \to \infty$,

$$\Pi_V(-q^2) = \frac{d_R}{24\pi^2} q^2 \ln(-q^2) + \cdots, \tag{15}$$

we match $g_5^2 = 12\pi^2/d_R$, where $d_R = N_{\mathrm{TC}}$ and $N_{\mathrm{TC}}(N_{\mathrm{TC}} + 1)/2$ are the dimensions for fundamental and symmetric second-rank tensor representations, respectively. We resort, however, to numerical analysis for $A(q, z)$ and correspondingly for the S parameter.

In technicolor theory, quarks and leptons get mass through coupling to techniquarks by an ETC interaction [15], whose scale has both lower and upper bounds, coming from the constraint to generate the observed mass, while suppressing the flavor-changing neutral-current processes. The QCD-like technicolor fails to generate an ETC scale that satisfies the constraint. However, if the β function of the technicolor theory has a quasi-IR fixed point, $\beta(\alpha_*) = 0$, the coupling is almost constant near the fixed point and thus the strength of the bilinear operator gets enhanced at low energy as

$$\bar{q}q|_{\Lambda_{\rm ETC}} \simeq \left(\frac{\Lambda_{\rm ETC}}{\Lambda_{\rm TC}}\right)^{\gamma_m} \bar{q}q|_{\Lambda_{\rm TC}}, \tag{16}$$

where γ_m is the anomalous dimension of the bilinear operator near the fixed point. The Schwinger-Dyson analysis shows $\gamma_m = 1$ in the phase where the chiral symmetry is broken [16,17]. Therefore, in the technicolor theories with a quasi-IR fixed point, which will be called conformal technicolor theories in short, a large enough mass for quarks and leptons is possible even with a large hierarchy between $\Lambda_{\rm ETC}$ and $\Lambda_{\rm TC}$.

The holographic dual of such conformal technicolor is very different from that of QCD-like technicolor. The mass of the bulk scalar becomes $m_5^2 = -4$, saturating the Breitenlohner-Freedman bound, since the scaling dimension of $\bar{q}q$ in the conformal technicolor is 2 instead of 3. The classical solution becomes $X_0(z) = c_1 z^2 + c_2 z^2 \ln(z/\epsilon)$. The UV boundary condition fixes the constant $c_1 = 0$ in the chiral limit and by AdS/CFT correspondence $c_2 = \langle\bar{q}q\rangle/4$.

To estimate the S parameter of conformal technicolor theories using holography, we solve

$$\left[z\partial_z\left(\frac{1}{z}\partial_z\right) + q^2 - \frac{g_5^2 X_0^2}{z^2}\right]A(q,z) = 0, \tag{17}$$

with the boundary conditions $A(q,\epsilon) = 1$ and $\partial_z A(q, z_m) = 0$. Since Eq. (17) is linear in $A(q,z)$ and $\Pi_A(q^2)$ is related to $\partial_z A(q,\epsilon)$, it is convenient to work instead with $T(q,z) \equiv \partial_z \ln A(q,z)$ with a *single* boundary condition, $T(q, z_m) = 0$. Then, $\Pi_A(-q^2)$ simply becomes $T(q,\epsilon)/(g_5^2\epsilon)$. Expanding $T(q,z) = T^{(0)}(z) + q^2 T^{(1)}(z) + \cdots$, we rewrite Eq. (17) as

$$z\partial_z\left(\frac{1}{z}T^{(0)}\right) + (T^{(0)})^2 = \frac{g_5^2 X_0^2}{z^2}, \tag{18}$$

$$z\partial_z\left(\frac{1}{z}T^{(1)}\right) + 2T^{(0)}T^{(1)} = -1, \tag{19}$$

and so on. Solving Eq. (19), we get

$$T^{(1)}(z) = -z\int_{z_m}^{z}\frac{dz'}{z'}\exp\left[2\int_{z}^{z'} d\omega T^{(0)}(\omega)\right]. \tag{20}$$

Since $g_5^2 \Pi_V(-q^2) = \ln(z_m/\epsilon)q^2 + O(q^4)$ for small q^2, the S parameter becomes

$$S = \frac{4\pi}{g_5^2}\int_{\epsilon}^{z_m}\frac{dz'}{z'}\left[1 - e^{2\int_{\epsilon}^{z'} d\omega T^{(0)}(\omega)}\right]. \tag{21}$$

To calculate the S parameter, we solve $T^{(0)}$ in (18) numerically with the boundary conditions $T^{(0)}(z_m) = 0$ and

$$-\Pi_A(0) = -\frac{1}{g_5^2}\frac{T^{(0)}(z)}{z}\bigg|_{\epsilon} = F_{\rm T}^2, \tag{22}$$

where the technipion decay constant $F_{\rm T} = 246\sqrt{2/N_{\rm TF}}$ GeV.

We solve Eq. (18) for the given boundary conditions, $T^{(0)}(z_m) = 0$ and $T^{(0)}(\epsilon)/\epsilon = -g_5^2 F_{\rm T}^2$, using two different bulk scalar fields: $X_0 = c_W z^2 \ln(z/\epsilon)$ for the technicolor with walking behavior (namely, for the techni-orientifold and the walking technicolor) and $X_0 = c_Q(z^3 - z\epsilon^2)$ for the QCD-like technicolor. Since $\partial_z T^{(0)}|_{z_m} = g_5^2 X_0^2/z_m^2$, there is a unique value for $c_{W,Q}$ or $\langle\bar{q}q\rangle$ which allows a solution for $T^{(0)}$. Therefore, in the holographic dual F_T determines $\langle\bar{q}q\rangle$ or vice versa.

In the ladder approximation the fermion self energy becomes $\Sigma(p) = \kappa/p$ for a large Euclidean momentum, $p \to \infty$, which then gives $\langle\bar{q}q\rangle = \Lambda\kappa$ at a scale Λ, if compared with the operator product expansion [16]. By analyzing the gap equation in the ladder approximation, one further finds that $\kappa = \Sigma^2(0)$ [17]. Since the dynamical mass of techniquarks $\Sigma(0) \simeq \Lambda_{\rm TC}$, $\langle\bar{q}q\rangle \simeq 1/z_m^3$ at $\Lambda_{\rm TC}$ or $c_W z_m^3 \simeq 1/4$.

Our results for the S parameter depend only on two dimensionless parameters, ϵ/z_m and $F_{\rm T} z_m$. For the former, we take $1/300 = \Lambda_{\rm TC}/\Lambda_{\rm ETC}$, while the latter must be estimated by other means. If we use the ladder approximation value $c_W z_m^3 = 1/4$, we obtain $F_{\rm T} z_m \simeq 0.86/g_5$ by solving Eq. (18). The corresponding S parameters are listed in the first row of Table I. If the chiral perturbation theory is valid up to the scale m_ρ, which is the first pole of Π_V in Eq. (14), we have $4\pi F_{\rm T} \simeq m_\rho$ or $F_{\rm T} \simeq 0.19/z_m$. Finally,

TABLE I. The S parameter for conformal technicolor with techniquarks in the symmetric second-rank tensor (S) and fundamental (F) representations.

$F_{\rm T} z_m$	$N = 2$, S	$N = 2$, F	$N = 3$, S	$N = 3$, F	$N = 4$, S	$N = 4$, F
$0.86/g_5$	0.086	0.057	0.17	0.086	0.29	0.12
0.19	0.15	0.14	0.17	0.15	0.17	0.16
0.29	0.28	0.22	0.34	0.26	0.37	0.31

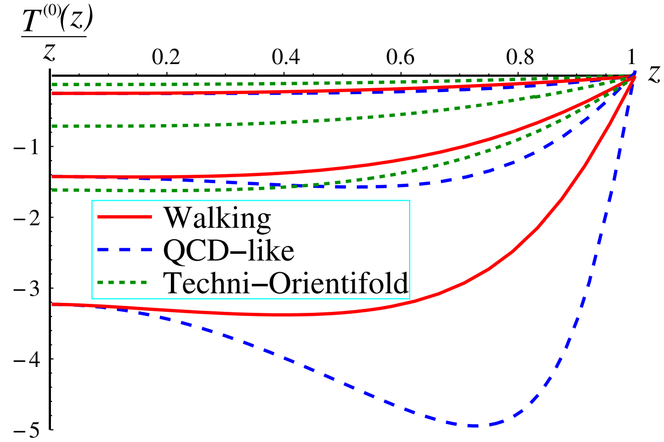

FIG. 1 (color online). The profiles for $N_{\rm TC} = 3$ with various $T^{(0)}(\epsilon)/\epsilon = -g_5^2 F_T^2$.

we include the QCD value, $m_\rho/F_{\rm T} \simeq 8.85$ or $F_{\rm T} \simeq 0.29/z_m$, for comparison. (For a strict comparison one should take into account that both $F_{\rm T}$ and z_m depend on $N_{\rm TC}$ and the representation of techniquarks.)

For the QCD-like technicolor with $F_{\rm T} = 0.29/z_m$, we find $S = 0.24$, 0.30, 0.34 for $N_{\rm TC} = 2$, 3, 4, respectively, which agrees well with the estimate by rescaling the QCD data [4]. In Table I, we list our numerical results for other technicolor models. Several comments are in order. First, we note that all the S parameters we calculated are positive. In fact, one can prove that the S parameter is always positive in the holographic dual of any models for dynamical electroweak symmetry breaking. By examining Eq. (18) one can show that $T^{(0)}$ is always negative for any value of X_0. Suppose $T^{(0)}(z)$ has one more zero for $z < z_m$. Since the right-hand side of Eq. (18) is always positive, the slope of $T^{(0)}/z$ has to be positive at those zeros, which is, however, impossible for a continuous function. Therefore $T^{(0)}$ must have no zeros and thus be negative for $z < z_m$. The S parameter in Eq. (21) is hence always positive.

Second, we find that the S parameter is reduced by about 10%–20% by walking when $F_{\rm T}$ and d_R are the same, which agrees with the Weinberg sum rule [18]. This can be seen easily, since the profile $-T^{(0)}(z)$ is always smaller for the walking technicolor than for the QCD-like technicolor with the same $F_{\rm T}$ and d_R and so is the S parameter. (See Fig. 1, where we take $z_m = 1$.) Suppose $\partial_z(T^{(0)}/z) = 0$ at $z = z_* > \epsilon$. Then, $\partial_z(T^{(0)}/z) \simeq h(z - z_*)$ near z_*. Expanding Eq. (18) around $z = z_*$ for the walking case, we find $h > 0$ and $z_* = \epsilon e^{h/(2g_5^2 c_W^2)}$. The slope of $T^{(0)}/z$ therefore vanishes only at a point very close to the UV boundary. $\partial_z(T^{(0)}/z)$ is positive for $0 < z \le z_m$ if we take $\epsilon \to 0$. On the other hand, for the QCD-like technicolor, the slope vanishes whenever (taking $\epsilon \to 0$)

$$\frac{1}{z} T^{(0)}(z) = -g_5 c_Q z. \tag{23}$$

Since $T^{(0)}(z)/z$ is finite at $z = 0$ but vanishes at z_m while the right-hand side of Eq. (23) is monotonically decreasing, there must exist a solution z_* to Eq. (23), remaining finite when $\epsilon \to 0$. We now note that the bulk fields X_0's for the walking technicolor and the QCD-like technicolor are monotonically increasing functions and meet together only at two points, $z = \epsilon$ and $z = z_0$. If $z_0 > z_m$, the slope of $T^{(0)}/z$ at z_m is bigger for the walking case and the $T^{(0)}/z$ profiles must meet at a point between z_* of the QCD-like technicolor and z_m, which is impossible, however, because then the slope has to be smaller for the walking case at the point, though its X_0 is bigger. Therefore $z_0 < z_m$ or the slope of $T^{(0)}/z$ at z_m has to be smaller for the walking technicolor and thus its profile $-T^{(0)}(z)$ is always smaller than that of QCD-like technicolor. The S parameter is therefore somewhat reduced by walking. However, we expect further reduction in the S parameter for the walking case, since $F_{\rm T}/\Lambda_{\rm TC}$ may be quite small in the walking technicolor [19].

Finally, we find that a slight change in $F_{\rm T}$ results in a substantial change in the S parameter. As shown in Fig. 1, the area sustained by $-T^{(0)}/z$ mainly depends on its value at $z = \epsilon$. If $F_{\rm T}$ is much smaller than $\Lambda_{\rm TC}$, the S parameter gets reduced substantially.

To conclude, we have analyzed the oblique S parameter in the holographic dual of technicolor theories. The AdS/CFT correspondence allows us to investigate the general aspects of the S parameter and also to obtain its numerical values precisely, which therefore removes a major hurdle for technicolor theories. We have also shown that the S parameter is strictly positive in the holographic dual models of technicolor and is reduced at least 10%–20% by walking. We predict $F_{\rm T}/\Lambda_{\rm TC} = 0.86/g_5$ in terms of the 5D gauge coupling in the ladder approximation, which results in small S parameters. For the technicolor models with $N_{\rm TF} > 2$, the S parameter will increase but not much since $F_{\rm T}$ decreases as $1/\sqrt{N_{\rm TF}}$.

The work of D. K. H. is supported by Korea Research Foundation Grant No. KRF-2005-015-C00108 funded by the Korean Government (MOEHRD), and H. U. Y. is supported by Grant No. R01-2003-000-10391-0 from the Basic Research Program of the Korea Science & Engineering Foundation.

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