1. Introduction

In the study of the H-mode [1], the bifurcation of radial electric field structure and subsequent suppression of cross-field transport [2–5] have played the role of central thread of thought to construct modeling (see, e.g. [6] for a review). Experimental evidences for the role of radial electric field were given [7, 8], and is summarized in review [9, 10]. Recent progress has provided the data of electric field structure near edge with high-spatial resolution (see, e.g. [11–15]). The dynamical response of barriers like self-organized limit cycle oscillation has attracted attentions [16–25], in collaboration with theoretical picture of L-H transitions [26–32]. The evolutions of the radial electric field and zonal flows [33], mean pressure gradient that induces radial electric field via neoclassical process, turbulence Reynolds stress, and others, have been measured in many devices.

In these achievements of the research, study of response of edge electric field structure against external biasing has provided a unique path to understand the nonlinear mechanism in bifurcation quantitatively [34–36]. A solitary radial electric field was observed [35, 36], and theoretical explanation has been discussed [37–39]. The induced bifurcation takes place under the condition where the spontaneous transition is difficult to occur. Thus this allows the study of nonlinearity in the neoclassical damping term, which is one of the origins of electric field bifurcation. In addition, even in the process of L-H transition, the magnitude of solitary electric field was found to jump without substantial change of ion pressure gradient [11]. This jump was discussed based on the electric field bifurcation model [40]. The physics of electric field bifurcation has been studied in helical plasmas as well [41–43].
radial electric field, which is induced by external biasing, was measured in LHD plasma [43]. Comparison of observations of tokamak and helical plasmas enriches understanding of toroidal plasma in general. Although there are many similarities in responses against biasing among tokamak and helical plasmas, a substantial difference in the magnitude of radial electric field is also noticeable. This stimulates the study of the scaling property of the solitary radial electric field, which is induced by the nonlinearity in the neoclassical damping process.

In this article, we extend the electric field bifurcation model in [38, 44] to the tokamak plasmas and helical plasmas, and study the dependence of the electric field structure on the plasma parameters and geometrical factors. The order of magnitude estimate for tokamak plasma is not far from experimental observations. It is shown that the height of electric field structure is reduced substantially owing to the ripple particle transport, while the width of the solitary electric field is influenced less. This might be useful for a basis to search the H-mode plasma with high confinement-enhancement factor in helical plasmas. In addition, the implication of the results to understand the limit of achievable gradient in the H-mode pedestal is also discussed.

2. Model

2.1. Model equation

The mean radial electric field $E_r$ is governed by the charge conservation relation combined with the Poisson’s relation, and is expressed as a nonlinear diffusion equation as (see chap.19 of [44])

$$\frac{\partial E_r}{\partial t} = \nabla \cdot \mu \nabla E_r - \frac{1}{\varepsilon_0 \varepsilon_1} (J_r - J_{\mathrm{ext}})$$

(1)

where $\mu$ is the ion viscosity,

$$\varepsilon_1 = 1 + \frac{e^2}{v_A^2} (1 + 2q^2)$$

(2)

is the dielectric constant in toroidal plasma (in plateau regime), $v_A$ is the Alfvén velocity, $q$ is the safety factor, $J_r$ is the current in the plasma, and $J_{\mathrm{ext}}$ is the component which is driven by external circuit. We are interested in the localized electric field structure, which is self-organized by the nonlinear response of the plasma. Thus the plasma parameters and their gradients (except $E_r$) in equation (1) are treated constant for the simplicity.

Normalization is introduced and the length, time, $E_r$ and current density are normalized as

$$x = \frac{(r - r_0)}{l} \quad \text{with} \quad l = \frac{\mu \varepsilon_0 \varepsilon_1}{\sigma(0)}, \quad (3a)$$

$$\tau = \frac{t}{t_a} \quad \text{with} \quad t_a = \frac{\varepsilon_0 \varepsilon_1}{\sigma(0)}, \quad (3b)$$

$$X = \frac{\rho_p E_r}{I}, \quad (3c)$$

where $\rho_p$ is the ion gyroradius at poloidal magnetic field, $\rho_p = q_i e^{-2} \rho$, $e$ is the inverse aspect ratio, $T_i$ is the ion temperature, and $\sigma(0)$ denotes a conductivity in the linear regime

$$\sigma(0) \equiv \partial J_r / \partial E_r |_{E_r \rightarrow 0}.$$ (3d)

The radius $r_0$ is the reference position, where the peak of solitary radial electric field appears, and is of the order of plasma minor radius. Then equation (1) is rewritten as

$$\frac{\partial}{\partial \tau} X = \frac{\partial^2}{\partial x^2} X - J(X) + I.$$ (4)

In this article, we study steady state solution, which is given by the equation

$$\frac{\partial^2}{\partial x^2} X - J(X) + I = 0.$$ (5)

In studying the stationary solution, the normalized external current $I$ is constant in space and time.

2.2. Solitary radial electric field structure

The solution of the solitary radial electric field is known to appear, when the current $J(X)$ has one maximum and the local equation

$$J(X) = I$$ (6)

has one stable branch and one unstable branch as is illustrated in figure 1. (In this article, the sign of electric field is chosen as positive in order to study the electric bias experiments. The extension to the negative electric field is straightforward.) The boundary condition is chosen as

$$X = X_A \quad \text{and} \quad dX/dx \to 0 \quad \text{at} \quad |x| \to \infty,$$ (7)

for the solitary radial electric field structure. An analytic solution for the radial electric field has been derived when $I$ is close to the peak value of the normalized current at $X = X_*$, as [44]

![Figure 1](image-url)
\[ X(x) = X_\lambda + (X_\rho - X_\lambda) \frac{12e^{\alpha x}}{(1 + e^{\alpha x})^2}, \]  
with 
\[ X(0) = X_\lambda + 3(X_\rho - X_\lambda). \]

Equation (8) denotes the solitary radial electric field (figure 2), which is localized near \( x = 0 \), where the parameter that characterizes the width of the peak is given as 
\[ \alpha = \sqrt{C(X_\rho - X_\lambda)}, \] 
with 
\[ C = -\frac{\partial^2 J(X)}{\partial X^2} \mid_{X=X_\lambda}. \] 
From equation (8a), one has an estimate, 
\[ X'' \sim -C(X_\rho - X_\lambda)^2. \] 

In this article, we focus upon the combination of the radial electric field and its curvature, \( XX'' \), as a key parameter for the suppression of turbulence [33]. A brief explanation is made in the appendix. It is evaluated as 
\[ XX'' \sim -CX_\rho(X_\rho - X_\lambda)^2. \]

### 3. Applications

#### 3.1. Bifurcation of electric field at tokamak edge

In tokamak configurations, the radial current by the neoclassical process has been derived as [37]
\[ J_r(E_r) = \frac{en\rho_p f}{BR^2} \frac{1}{\sqrt{\pi}} \text{Im}Z(X+i\omega \alpha^{-1})(E_r - E_{e,a}). \]

where \( Z(X+i\omega \alpha^{-1}) \) is the plasma dispersion function, the argument of which is \( X+i\omega \alpha^{-1} \), \( \nu_i \) is the ion-ion collision frequency, \( \alpha_i = \nu_i/qR \) is the ion transit angular frequency, \( E_{e,a} \) represents the radial electric field which is induced by the neoclassical bulk viscosity, and is normalized as,
\[ \chi_a = \rho_p \nu_i^{-1} e E_{e,a} = \rho_p (\alpha n^{-1} + \gamma_i T_i^{-1}) + V_i \nu_i^{-1}, \]
\( \gamma_i \) is the specific heat ratio of ions, and \( V_i \) is the mean ion velocity along the magnetic field. From equation (12a), one has the conductivity (in the limit of weak electric field) as
\[ \sigma(0) = \frac{en\rho_p f}{BR^2} \]
and the normalized current function as
\[ J(X) = \frac{1}{\sqrt{\pi}} \text{Im}Z(X+i\omega \alpha^{-1})(X - X_a). \]

We are interested in the case that the bifurcation to the solitary radial electric field structure occurs away from the spontaneous L-H condition, so that the simplification
\[ |X_a| \ll |X| \]
is employed. This condition (15) may also be used in the circumstance where the jump of radial electric field is large enough (in comparison with the ion-pressure-gradient driven radial electric field) in the H-mode plasma, as was reported in [11].

From equations (14) and (15), we have an order of magnitude estimate
\[ X_r \sim 1, \]
and
\[ \frac{\partial^2 J(X)}{\partial X^2} \mid_{X=X_r} \sim -1. \]

See, e.g. [37] for more accurate evaluation of the second derivative of \( J(X) \). Thus, one has
\[ XX'' \sim -1 \]

apart from a numerical factor of the order unity) in the normalized unit.

An order of magnitude estimate of the radial electric field is given in the experimental units, by rewriting equation (16) by use of equation (3c), as
\[ |E_r| \approx \frac{T_i}{e\rho_p}. \]

Equation (13) for the conductivity gives the ratio of \( \sigma(0) \) to dielectric constant as
\[ \frac{\sigma(0)}{e\epsilon_0} \approx \alpha_i. \]

Substituting this relation (20) into the formula of normalizing length, equation (3a), one has
\[ l^2 \approx \mu_i \alpha_i^{-1}. \]

This length \( l \) denotes the characteristic value for the width of the peaked profile. Following the similar process, equation (18) is expressed in the dimensional form as
\[ E_r \chi_a' \approx -\frac{\alpha_i}{\mu_i} \left( \frac{T_i}{e\rho_p} \right)^2. \]

Equations (21) and (22) show that the solitary structure becomes steeper when the turbulence transport is suppressed. If one uses a neoclassical value for the ion shear viscosity,
considering the situation where the turbulence transport is suppressed, as
\[ \mu_t \sim \chi_{NC} \sim O(\nu_\perp J_q^2 \rho_p^2) \]  \hspace{1cm} (23)
for the analytic insight of the problem, one has the characteristic radial scale length of the solitoy electric field from equations (21) and (23) as
\[ l \sim \nu_\perp J_q. \]  \hspace{1cm} (24a)
This is the estimate of the upper bound of steepness of the solution. The peak curvature of electric field, combined with the electric field strength, is estimated from equations (22) and (23) as
\[ -E_r E_r' \approx \frac{1}{\nu_\perp J_q^2 \rho_p^2} \left( \frac{T_i}{e \rho_p^2} \right)^2 \]  \hspace{1cm} (24b)
It is characterized by the ion temperature and poloidal gyroradius.

One might be interested in how the result behaves in the banana regime. In comparison with the plateau regime, the perpendicular dielectric constant and shear viscosity are modified as
\[ \varepsilon_i(\text{banana}) \sim \frac{1}{\sqrt{\nu_\perp J_q}} \varepsilon_i(\text{plateau}), \]  \hspace{1cm} (25a)
and
\[ \mu_i(\text{banana}) \sim \frac{\nu_\perp J_q^2 \omega_\perp}{\varepsilon_i(\text{plateau})}. \]  \hspace{1cm} (25b)
By use of these relations, equations (24a) and (24b) are rewritten as
\[ l \sim \frac{\nu_\perp J_q}{\omega_\perp \rho_p} \]  \hspace{1cm} (26a)
and
\[ -E_r E_r' \sim \frac{\omega_\perp}{\nu_\perp J_q^2 \rho_p^2} \left( \frac{T_i}{e \rho_p^2} \right)^2 \]  \hspace{1cm} (26b)
in the banana regime.

3.2. Case of helical plasmas
This model is applied to helical plasmas, in order to examine the observation in [43]. We take a simple model of magnetic field for the helical plasma, the multi polarity of which is 2, as
\[ B = B_0 (1 - e_h \cos(2\theta - M \zeta) - e_i \cos \theta), \]  \hspace{1cm} (27)
where \( e_h \) is the helical ripple amplitude, \( \theta \) and \( \zeta \) are the poloidal and toroidal angles, respectively, and \( M \) is the number of field period. The regime of helical ripple transport,
\[ \nu_{\rm eff,h} \equiv \nu_t/2e_h \ll \omega_h \equiv \sqrt{2e_h v_T MR^{-1}} \]  \hspace{1cm} (28a)
is chosen. The ratio between the effective collision frequency (LHS) and the transit frequency in the helical ripple (RHS) is sometimes referred to as
\[ \nu_{\rm \perp,h} \equiv \frac{\nu_{\rm eff,h}}{\omega_h} \]  \hspace{1cm} (28b)
In addition, we study the case of so-called ‘ion-root’, i.e. the helical ripple transport of ions in the presence of radial electric field is much larger than that of electrons,
\[ \Gamma(E_r = 0) \gg \Gamma(E_r = 0). \]
The modification of the ion-root branch by external current is studied here. The radial current is approximately given by
\[ J_r = e (T_i - \Gamma_e) \sim e \Gamma_e. \]  \hspace{1cm} (29)
This is for the clarification of the problem, and the study of general cases, in which change between the electron-root and ion-root can take place, is left for future work. The radial particle flux of ions has been derived in [45] as
\[ \Gamma_i = D_{00} \left( \frac{e E_r}{T_i} - \frac{c_i e T_i'}{n} - \frac{c_i T_i'}{T_i} \right) \]  \hspace{1cm} (30)
where \( c_a \) and \( c_T \) indicate the neoclassical coefficients, and the fitting formula of the ripple transport coefficient has been derived as
\[ D_i = \frac{D_{00}}{1 + S^{3/2} X^{3/2}}, \]  \hspace{1cm} (31a)
\[ D_{00} \approx \frac{247}{9 \pi} \frac{\varepsilon e BR}{\nu_{\rm \perp,h}^2} \]  \hspace{1cm} (31b)
and
\[ S \approx \frac{7.5 \nu_{\rm \perp,h}^3}{M q_i e BR} \]  \hspace{1cm} (31c)
which is related to equation (28) as
\[ S \approx \frac{7.5}{M q_i e BR} \frac{1}{\nu_{\rm \perp,h}}. \]  \hspace{1cm} (31d)
Comparing the condition for the ripple transport regime, equation (28),
\[ \nu_{\rm \perp,h} \ll 1, \]  \hspace{1cm} (32a)
one sees that the condition
\[ S \gg 1 \]  \hspace{1cm} (32b)
holds in the regime of helical ripple transport. Equations (31a) and (32b) indicate that the ripple particle transport is more strongly reduced by the radial electric field, in comparison with the case of tokamak plasmas. By employing the normalization of equation (3), one has
\[ \sigma(0) = \frac{\partial J_r}{\partial E_r} \bigg|_{E_r=0} \equiv \frac{e^2 n D_{00}}{T_i} \]  \hspace{1cm} (33)
and the normalized current function
\[ J(X) = \frac{1}{1 + S^{3/2} X^{3/2}} (X - X_{a,b}), \]  \hspace{1cm} (34a)
where
\[ X_{a,b} = \rho_p (c_i n^{-1} + c_T T_i T_i^{-1}) \]  \hspace{1cm} (34b)
is the contribution of neoclassical bulk viscosity effect, and coefficients are given as \( \hat{c}_n = c_\alpha (1 + S^{3/2} k^{3/2}) \) and \( \hat{c}_T = c_T (1 + S^{3/2} k^{3/2}) \). As is the case in the preceding section, we neglect the contribution of \( X_n \), and use a simplified form

\[
J(X) = \frac{X}{1 + S^{3/2} \lambda^{3/2}}. \tag{35}
\]

Model equation (35) tells that the peak of the current function appears at

\[
X_n \sim 1.6S^{-1} \tag{36}
\]

and it gives

\[
\frac{\partial^2 J(X)}{\partial X^2}|_{X=X_n} \sim -0.1S. \tag{37}
\]

Substituting equations (36) and (37) into equation (10), one has (apart from a numerical factor of the order unity)

\[
X^s \sim -S^{-1}. \tag{38a}
\]

With the help of equations (36) and (38a), the peak curvature of electric field, combined with the electric field strength, is estimated as,

\[
XX^s \sim -S^{-2} \tag{38b}
\]

in the normalized unit.

Equations (3a) and (33) give the normalization length as

\[
l^2 = \frac{\mu}{D_0} q_s^2 \lambda_s^2. \tag{39}
\]

Considering the case that the helical ripple transport is dominant in the transport process, one takes an approximation of

\[
\mu_l / D_0 \sim 1. \tag{40}
\]

Substituting equation (40) into equation (39), one has

\[
l \sim \epsilon_0 \rho_p \tag{41}
\]

in the experimental unit. Then, equation (38b) is rewritten in the dimensional unit as

\[
-E_s E_p^s \simeq \frac{1}{S^2} \frac{1}{\epsilon^2 \rho_p^2} \left( \frac{T_e}{\epsilon_0 \rho_p} \right)^2. \tag{42}
\]

Comparing results for tokamaks and helical systems, equations (24b) and (42), respectively, one sees that the electric field curvature effect is reduced by the factor \( S^{-2} \) if helical ripples affect the transport. The difference between helical systems and tokamaks appears in the magnitude of the localized electric field, and the radial scale length obeys similar dependence.

### 4. Summary and discussions

In this article, the solitary radial electric field structure in the edge of toroidal plasmas are investigated. Based on the electric field bifurcation model, order of magnitude estimates of the electric field, curvature and radial scale length are given. The result is summarized in the table 1. It is noted that the scale length and curvature depend on the model of the shear viscosity of ions, but the magnitude of peak electric field is independent of the model of shear viscosity.

Although the results are obtained as scaling relations (apart from numerical coefficients of the order of unity), it might be worth examining the absolute values which the model provides. For the parameters of the tokamak edge (see, [13] for detailed measurements), equation (24b) or (26) gives a value of the order of a few \( 10^{15} [V^2 m^{-1}] \) for the peak curvature of electric field multiplied by the electric field strength. Experiments have reported the number of about \( 10^{15} [V^2 m^{-1}] \) [13, 15, 36]. Thus the scaling relations (table 1) is not bad in understanding the experimentally observed solitary radial electric field structure in tokamaks. For the case of helical plasmas, the value (for the peak curvature of electric field multiplied by the electric field strength) is suppressed by the factor of \( S^{-2} \) owing to the helical-ripple trapped particles, while the width of the peak is independent of \( S \). Considering the fact that the parameter \( S \) is in the range of a few \( 10^4 \), the reduction factor is about \( \sim 10^3 \) times smaller compared to tokamaks. Experiments have reported the value of \( \sim 10^9 [V^2 m^{-1}] \) [43]. The width of the solitary peak is not far from those in tokamaks. Thus, the scaling relations in this article give qualitative understanding for the differences in the solitary radial electric field structures between tokamaks and helical plasmas.

When the electric field inhomogeneity is weak and turbulence is not suppressed, the viscosity is larger than equation (23). If one substitutes the gyro-Bohm dependence, \( \mu_l \sim D_{\nu B} \sim O(\omega_0 \rho_p \rho), \) one has a dependence \( l \approx \sqrt{\epsilon_0 / \epsilon \rho_p}. \) The curvature of the electric field (multiplied by electric field) is modified accordingly as, \(-E_s E_p^s \sim \epsilon^{-1} \rho_p^2 (T_e / \epsilon_0 \rho_p)^2 \), while the scaling relation of the peak value of radial electric field is not influenced.

It should be noticed that, the scaling relations in table 1 is obtained as an order of magnitude estimate. One must examine by the numerical calculation whether the width of radial electric field is of the order of or smaller than the banana width. In order to obtain the accurate number, one must employ more relevant formula of transport coefficient (in which the squeezing of the banana particle orbits is taken into account). Such an extension is left for future study.

The term \( XX^s \) has been predicted important in suppressing turbulence via modulational coupling [33]. The suppression is prominent if the rate of modulation of fluctuations by the

|                | \(-E_s E_p^s\) | \(E_s\) | \(l\) |
|----------------|----------------|-------|-----|
| Tokamak (plateau regime) | \(1 / \epsilon^2 \rho_p^2 \left( T_e / \epsilon_0 \rho_p \right)^2 \) | \(T_e / \epsilon_0 \rho_p\) | \(\epsilon_0 \rho_p\) |
| Tokamak (banana regime) | \(a_t / \epsilon_0 \rho_p^2 \left( T_e / \epsilon_0 \rho_p \right)^2 \) | \(T_e / \epsilon_0 \rho_p\) | \(\sqrt{a_t / \epsilon_0 \rho_p}\) |
| Helical | \(1 / S^2 \epsilon^2 \rho_p^2 \left( T_e / \epsilon_0 \rho_p \right)^2 \) | \(1 / S T_e / \epsilon_0 \rho_p\) | \(\epsilon_0 \rho_p\) |
electric field is of the order of the decorrelation rate of microscopic drift wave turbulence. The nonlinear decorrelation rate of drift wave fluctuations is evaluated as

$$\omega \sim c_s / L_n$$

for $k_d p_1 \sim 1$, where $k_d$ is the wavenumber of drift wave fluctuation and $L_n$ is the density gradient scale length. The condition that the rate of modulation by the inhomogeneous radial electric field is larger than the nonlinear decorrelation rate of fluctuations,

$$\sqrt{\frac{E_r E_r'}{B^2}} > \omega \sim c_s / L_n,$$  \hfill (44)

is rewritten as

$$\frac{1}{q} > \frac{\rho_p}{q_n},$$ \hfill (45)

for tokamaks. Here, equation (24b) is used for plateau regime. For given geometrical parameters, the suppression of turbulence (by the bifurcation to the state with solitary radial electric field structure) stops when the gradiel becomes steep and the gradient scale length becomes of the order of $q_p / \rho_p$. This describes a scaling property for one of the achievable limits of density gradient in the H-mode pedestal. For helical systems, combination of equations (42) and (44) yields the scaling relation for the condition as

$$\frac{1}{S q} > \frac{\rho_p}{q_n},$$ \hfill (46)

in the helical ripple particle regime. This condition (46) is more stringent than equation (45). The achievable gradient is weaker by the factor $S^{-1}$ than equation (45). Thus, in helical plasmas, the effective suppression of turbulence by the solitary radial electric field occurs more easily in the plateau regime than in helical ripple particle regime, if other parameters are common.

The impacts of magnetic field ripples on the edge barrier in tokamaks have attracted attentions, particularly in conjunction with symmetry-breaking-magnetic perturbation for ELM control [47]. As has been discussed in [40, 48], the influence of small ripples is complicated, and may introduce a new bifurcation in addition to the one, which is discussed in this article. The issue requires future intensive studies.

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Appendix A. Curvature of radial electric field

We here briefly explain that the curvature of radial electric field is focused upon in considering the suppression of microscopic fluctuations via modulational coupling. When the turbulence Reynolds stress is induced by the microscopic fluctuations via disparate-scale interaction, the turbulent Reynolds stress is proportional to the gradient of the radial electric field, i.e.

$$\langle \tilde{v} r \tilde{v}_r \rangle \propto \frac{d}{dr} E_r,$$ \hfill (A.1)

where $\tilde{v}_r$ is the fluctuation velocity by microscopic fluctuations [33]. This is natural from symmetry consideration. Thus, the force by this Reynolds stress per unit mass density, $d\langle \tilde{v} r \tilde{v}_r \rangle / dr$, is proportional to the curvature of the radial electric field. The power absorbed from turbulence by the flow is proportional to the force multiplied by velocity, $V_{Esb} d\langle \tilde{v} r \tilde{v}_r \rangle / dr$, and has a relation

$$V_{Esb} \frac{d}{dr} \langle \tilde{v} r \tilde{v}_r \rangle \propto -E_r \frac{d^2}{dr^2} E_r,$$ \hfill (A.2)

Therefore, the product of electric field and its curvature, \(XX''\) in normalized variable, has a key role in the suppression of turbulence via modulational coupling. This is also seen from the fact that the growth rate of the zonal flow is proportional to $q^2$, i.e. its curvature, as is shown by equation (3.2.21) of [33] (where $q_r$ is the wave number of the zonal flow). The turbulence intensity is sensitive to the curvature of flow.

The relation with the argument on the electric field shear is as follows. In the 0D model, quantities are averaged over the localization length. In spatially-averaging the energy exchange rate equation (A.2) over the localization length, the partial integral gives

$$- \int dr E_r \frac{d^2}{dr^2} E_r = \int dr |E_r|^2 = \langle U \rangle,$$ \hfill (A.3)

where $U$ is the vorticity of $E \times B$ flow. Thus, the electric field shear is often referred to, for an order of magnitude estimate of the process equation (A.2). Other mechanism that the electric field shear can suppress the turbulence is stretching of eddies. As is explained in [3], the stretching of turbulence eddies by the electric field shear enhances the decorrelation of drift wave fluctuations. That is, this process does not transfer the energy of fluctuations to the mean electric field. Therefore, there are (at least) two independent processes due to the inhomogeneous electric field, by which fluctuations loose energy. The detailed measurement of the edge transport barrier on JT-60U has shown that the peak of gradient at barrier appears at the peak of curvature of radial electric field, not at the peak of gradient of radial electric field [11–13]. The observation on LHD has also shown that the fluctuations are strongly influenced by electric field curvature [43]. It suggests that, under such circumstances, the influence of curvature is stronger than that of shear. Thus, we investigate the quantity \(XX''\) in this manuscript.

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