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Subcritical Excitation of Plasma Turbulence

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Abstract

Theory of current-diffusive interchange mode turbulence in plasmas is developed in the presence of collisional transport. Double-valued amplitude of stationary fluctuations is expressed in terms of the pressure gradient. The backward bifurcation is shown to appear near the linear stability boundary. The subcritical nature of the turbulence is explicitly illustrated. Critical pressure gradient at which the transition from collisional transport to the turbulent one is to occur is predicted. This provides a prototype of the transport theory for nonlinear-non-equilibrium systems.

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Keywords: subcritical turbulence, interchange mode, current diffusivity, renormalization, dressed test mode, critical pressure gradient

1. Introduction

Fluctuations and associated transport are important issues in the physics of non-equilibrium systems. Anomalous transport phenomenon in confined plasmas is one of typical examples, and has been subject to intensive study [1]. Theory for the transport due to the binary collision of particles has been established [2], nevertheless it is insufficient to explain many of the plasma dynamics. To investigate anomalous transport phenomena in confined plasmas, much work has been done on the linear stability of microscopic modes; to obtain the saturation level, the balance between the linear growth and nonlinear damping has been discussed [3]. Nonlinear damping mechanisms are demonstrated in numerical simulations [4], but recent simulations have shown that the plasma turbulence can belong to a class of subcritical turbulence [5,6]. The subcritical turbulence has been studied in fluid dynamics, and the conventional methods like amplitude expansion or truncation of modes often encounter the difficulty of the poor convergence [7]. To analyze the subcritical turbulence, analysis nearby the equilibrium is not sufficient. The theory of subcritical turbulence in plasmas must be advanced in order to understand the fluctuations and anomalous transport in confined plasmas.

Recently, a theoretical method of self-sustained turbulence was proposed in the analysis of magnetically-confined plasmas [8]. In this method, the $v \nabla v$ nonlinearity is renormalized in a form of the turbulent-driven diffusivities. Dynamical equation for the "dressed-test mode" was derived. In this paper, analysis using the dressed test mode is made in the presence of both the thermal fluctuations (i.e., the Coulombic collisions in magnetized plasmas) and the turbulence. The interchange mode turbulence, which is the simplest example in the presence of inhomogeneous magnetic and plasma pressures, is analyzed. The fluctuation amplitude is *analytically* expressed in terms of the global pressure gradient, which clearly shows the feature of the subcritical turbulence. Critical pressure gradient at which the transition from the collisional transport to the anomalous one is to occur is obtained. Comparison study with numerical simulation is made and shows comprehensive agreement.

2. Model Equation and Renormalization

We study the high-aspect-ratio, toroidal helical plasma with magnetic hill and strong magnetic shear. The minor and major radii of the torus are given by a and R, respectively, We use the toroidal coordinate (r, θ, ζ). The reduced set of equations for the electrostatic potential ϕ, pressure p, and current J are employed [9]. Equation of motion: $\frac{\partial \nabla_\perp^2 \phi}{\partial t} + [\phi, \nabla_\perp^2 \phi] = \nabla_\parallel J + (\Omega' \times \hat{\zeta}) \cdot \nabla p + \mu_{\perp c} \nabla_\perp^4 \phi$, the Ohm's law: $\frac{\partial \Psi}{\partial t} = -\nabla_\parallel \phi - \frac{1}{\xi}\left(\frac{\partial J}{\partial t} + [\phi, J] \right) + \lambda_c \nabla_\perp^2 J$ and the energy balance equation: $\frac{\partial p}{\partial t} + [\phi, p] = \chi_c \nabla_\perp^2 p$. The bracket [f,g] is the Poisson bracket, $[f, g] = (\nabla f \times \nabla g) \cdot \vec{b}$, ($\vec{b} = \vec{B}_0 / B_0$, B_0 being the main magnetic field), Ω' is the average curvature of the magnetic field, Ψ is the vector potential, and $1/\xi$ denotes the finite electron inertia, $1/\xi = (\delta/a)^2$, δ being the collisionless skin depth. The classical resistivity is neglected. The transport coefficients $\mu_{\perp c}$, λ_c, χ_c are the contributions from collisional diffusion and are the viscosity for the perpendicular momentum, the current diffusivity and the thermal diffusivity, respectively. In this article length and time are normalized to a and poloidal Alfven transit time τ_{Ap}. Pressure and potential are normalized to $B_0^2/2\varepsilon\mu_0$ and $a v_{Ap} B_0$, ($\varepsilon = a/R$). This set of equations is very much simplified in the plasma geometry, but is relevant to study the anomalous transport in the system of magnetic hill such as torsatron/Heliotron [10] and the inside of q=1 surface of tokamaks.

The Lagrangian nonlinearity is included in a form of $[\phi, \cdots]$ in the set of basic equations. These terms are renormalized as follows. (Detailed procedure is given in [11].) We take a test mode (denoted by k) and study the interaction with back-ground fluctuations (denoted by k_1). The driven mode (k_2), which is generated by the mode coupling between the test mode and fluctuations, are calculated. The back-interaction of the driven mode with back-ground fluctuations generates the test mode. Taking this nonlinear process (called direct interactions), the nonlinear term for the test mode, $[\phi, \tilde{Y}]$, is written as $[\phi, \tilde{Y}] \cong \sum_{k_1} [\phi_{-1}, [\phi_1, \tilde{Y}]]$, where $\tilde{Y}$ is a component of the test mode. (The suffix 1 indicates k_1.) The spatial inhomogeneity of fluctuations is assumed to be much weaker than the mode number, and the diffusion approximation is employed as

$[\phi_{-1}, [\phi_1, \tilde{Y}]] = (|\partial\phi_1/\partial r|^2 \partial^2\tilde{Y}/r^2\partial\theta^2 + |\partial\phi_1/r\partial\theta|^2 \partial^2\tilde{Y}/\partial r^2)$. Without loosing generality, the simplification of the isotropic turbulence, $\langle|\partial\phi_1/\partial r|^2\rangle \simeq \langle|\partial\phi_1/r\partial\theta|^2\rangle \simeq \langle|k_{\perp 1}\phi_1|^2\rangle/2$, is made. Bracket $\langle\rangle$ means a spectrum average. Only the diagonal elements are kept in the following. After these renormalization processes, the set of basic equations is expressed in terms of the dressed test mode with renormalized transport coefficients as

$$\frac{\partial\nabla_\perp^2\phi}{\partial t} = \nabla_\parallel J + (\Omega'\times\hat{\xi})\cdot\nabla p + (\mu_{\perp N} + \mu_{\perp c})\nabla_\perp^4\phi \tag{1}$$

$$\frac{\partial\Psi}{\partial t} = -\nabla_\parallel \phi - \frac{1}{\xi}\frac{\partial J}{\partial t} + (\lambda_N + \lambda_c)\nabla_\perp^2 J \tag{2}$$

$$\frac{\partial p}{\partial t} = (\chi_N + \chi_c)\nabla_\perp^2 p - [\phi, p_0] \tag{3}$$

In deriving these equations, the nonlinear diffusion coefficients for the test mode, $\{\mu_{\perp k}, \lambda_k, \chi_k\}$, due to the back-ground turbulence appears, which in principle can depend on the choice of k. Then the mean field approximation is employed: the range of the test mode k is not distinguished from that of the back-ground turbulence, and common ranges are taken for k_1 and k_2. Summing up over the background fluctuations, k_1, a set of diffusion coefficients $\{\mu_{\perp N}, \lambda_N, \chi_N\}$ was explicitly given as $\mu_N = \Sigma \frac{|k_{\perp 1}\varphi_1|^2}{2}\frac{1}{K_1}$, $\mu_{eN} = \Sigma \frac{|k_{\perp 1}\varphi_1|^2}{2}\frac{1}{K_1}\frac{\gamma_{u1}}{\gamma_{j1}}\left[1 + \frac{k_{\theta 1}^2 G_0}{\gamma_{u1}\gamma_{p1}k_{\perp 1}^2}\right]$,
$\chi_N = \Sigma \frac{|k_{\perp 1}\varphi_1|^2}{2}\frac{1}{K_1}\frac{\gamma_{u1}}{\gamma_{p1}}\left[1 + \frac{\xi k_{\parallel 1}^2}{k_{\perp 1}^2\gamma_{u1}\gamma_{j1}}\right]$, $K_1 = \gamma_{u1} + \frac{k_{\parallel 1}^2}{k_{\perp 1}^2}\frac{\xi}{\gamma_{j1}} + \frac{k_{\theta 1}^2 G_0}{\gamma_{p1}k_{\perp 1}^2}$ and $\lambda = (\delta/a)^2\mu_e$, where $\gamma_{u1} = \gamma(1) + \Gamma_{u1}$, $\gamma_{j1} = \gamma(1) + \Gamma_{j1}$, $\gamma_{p1} = \gamma(1) + \Gamma_{p1}$ $\gamma(1)$ is the eigenvalue of the k_1 mode, $\partial\{U_1, J_1, p_1\}/\partial t = \gamma(1)\{U_1, J_1, p_1\}$, Γ_{u1}, Γ_{j1}, and Γ_{p1} denote the decorrelation rate of U_1, J_1, and p_1. Other notation is: U is the vorticity, $U = -k_\perp^2\phi$, G_0 is the normalized pressure gradient, $G_0 = \Omega'\frac{dp_0}{dr}$, and p_0 is the equilibrium pressure profile. The transport coefficients, which operate to the test mode, are explicitly expressed in the presence of the collisional components μ_c, λ_c, χ_c and turbulent terms μ_N, λ_N, χ_N .

3. Stationary Solution and Backward Bifurcation

The stationary solution is studied. The limit of $\gamma \to 0$ is taken, and the relations $\gamma_{u1} \approx (\mu_N + \mu_c)k_{1\perp}^2$, $\gamma_{j1} \approx (\mu_{eN} + \mu_{ec})k_{1\perp}^2$, $\gamma_{p1} \approx (\chi_N + \chi_c)k_{1\perp}^2$ and $K_1 = (\mu_N + \mu_c)k_{1\perp}^2(1 + C)$ are obtained, where C is defined by $C = \xi k_{\parallel 1}^2 k_{\perp 1}^{-2}\gamma_{u1}^{-1}\gamma_{j1}^{-1} + G_0 k_{\theta 1}^2 k_{\perp 1}^{-2}\gamma_{u1}^{-1}\gamma_{p1}^{-1}$. (Note that λ is related to the electron viscosity as $\lambda_{(N,c)} = (c/a\omega_p)^2\mu_{e(N,c)}$.) In this limit, we have the relation between the turbulent transport coefficient and fluctuation amplitude as

$$\mu_N(\mu_N + \mu_c) = \hat{\phi}^2 \tag{4-1}$$

$$\mu_{eN}(\mu_{eN} + \mu_{ec}) = P^2\hat{\phi}^2 \tag{4-2}$$

$$\chi_N(\chi_N + \chi_c) = Q^2\hat{\phi}^2 \tag{4-3}$$

where the normalized fluctuation amplitude $\hat{\phi}$ is defined as $\hat{\phi}^2 \equiv \Sigma|\varphi_1|^2(2 + 2C)^{-1}$. The coefficients P and Q, (which are related to Prandtl numbers) are defined as

$$P^2 \equiv \left\langle \Sigma|\varphi_1|^2\left[1 + \frac{k_{\theta 1}^2 G_0}{\gamma_{u1}\gamma_{p1}k_{\perp 1}^2}\right](1 + C)^{-1}\right\rangle \left\{\Sigma|\varphi_1|^2(1 + C)^{-1}\right\}^{-1}$$

$$Q^2 \equiv \left\langle \Sigma|\varphi_1|^2\left[1 + \frac{\xi k_{\parallel 1}^2}{k_{\perp 1}^2\gamma_{u1}\gamma_{j1}}\right](1 + C)^{-1}\right\rangle \left\{\Sigma|\varphi_1|^2(1 + C)^{-1}\right\}^{-1}$$

The coefficients P and Q are the ratios of the moments of turbulent fluctuation spectrum, and vary much more slowly than the turbulence level itself. They are approximated as constants. In the strong turbulent limit, ratios $\mu_{eN}/\mu_N = P$ and $\chi_N/\mu_N = Q$ are found to be close to unity [12].

Basic equations (1)-(3) are linearlized for the dressed test mode, so that the nonlinear marginal stability condition is derived. (See [6,11] for the details of the solution of the eigenmode equation.) The marginal stability condition for the least stable mode was derived as

$$\frac{G_0^{3/2}}{s^2}\frac{(\lambda_N+\lambda_c)}{(\chi_N+\chi_c)^{1.5}(\mu_N+\mu_c)^{0.5}}=\Im_c \qquad (5)$$

where $\Im_c$ is a critical Itoh-number and is of the order of unity. Equations (4) and (5) determine the fluctuation level and turbulent transport coefficient as a function of the equilibrium pressure gradient, i.e., $\hat{\phi}(G_0)$ and $\chi(G_0)$.

In order to examine the bifurcation nature of this mode, we consider the neutral condition in the vicinity of the linear boundary. The linear stability is obtained by taking the limit of $\hat{\phi}\to 0$ as

$$G_0 \le G_c \qquad (6)$$

where the critical pressure gradient is given from Eq.(5) with $\mu_N=\lambda_N=\chi_N=0$ as $G_c=\Im_c^{2/3}(sa\omega_p/c)^{4/3}\chi_c\mu_c^{1/3}\chi_{ec}^{-2/3}$. By expanding Eq.(5) near $G_0\simeq G_c$ as $G_0\simeq G_c+(\partial G_0/\partial\hat{\phi}^2)\hat{\phi}^2+\cdots$, the amplitude ϕ near the marginal condition is given as $\hat{\phi}^2=(\partial G_0/\partial\hat{\phi}^2)^{-1}(G_0-G_c)$. From Eq.(5), the derivative is calculated as $\partial G_0/\partial\hat{\phi}^2=\{(\partial\chi_N/\partial\hat{\phi}^2)/(\chi_c+\chi_N)+(\partial\mu_N/\partial\hat{\phi}^2)/3(\mu_c+\mu_N)-2(\partial\mu_{eN}/\partial\hat{\phi}^2)/3(\mu_{ec}+\mu_{eN})\}G_0$. Noting the relation Eq.(4) and taking the limit of $\hat{\phi}\to 0$, we have $\partial G_0/\partial\hat{\phi}^2=(Q^2/\chi_c^2+1/3\mu_c^2-2P^2/3\mu_{ec}^2)G_c$ near $G_0\simeq G_c$. Relations $\mu_c\simeq\chi_c$ and $\mu_{ec}\ll\chi_c$ usually hold. For collisional diffusion, the relation $\mu_{e,c}/\chi_c\sim\sqrt{m_e/m_i}$ holds, where m_e/m_i is the mass ratio. Under this circumstance, the relation $\partial G_0/\partial\hat{\phi}^2<0$ holds near $G_0\simeq G_c$, and

$$\hat{\phi}^2=-(3\mu_{ec}^2/2P^2)(G_0/G_c-1) \qquad (7)$$

This result shows the backward-bifurcation: small but finite amplitude of ϕ is expected at the pressure gradient that is below the critical pressure gradient, $G_0<G_c$.

The branch of the large fluctuation amplitude is also obtained from Eqs. (4) and (5). In the large amplitude limit, $\chi_N\gg\chi_c$, Eqs.(4) and (5) reduce to the relation

$$\chi = \hat{\phi} = \frac{G_0^{3/2}}{s^2}\left(\frac{c}{a\omega_p}\right)^2 \qquad (8)$$

where the typical mode number scales as $k_\perp \sim G_0^{-1/2} sa\omega_p/c$. This branch corresponds to the strong turbulence limit of the self-sustained turbulence. The lower branch (7) and upper branch (8) are found to merge at particular pressure gradient, $G_0 = G_*$. The merging point is given from the singularity condition, $\partial G_0/\partial\hat{\phi} = 0$. This condition is satisfied at $\chi_N \simeq \chi_c - \mu_{e\,c}$, in the limit of $\chi_c \sim \mu_c >> \mu_{e\,c}$, with the pressure gradient

$$G_0 = G_* \cong \Im_c^{2/3}(2sa\omega_p/c)^{4/3}\chi_c^{2/3} \qquad (9)$$

At the critical pressure gradient G_*, the turbulent transport coefficient is expected to be of the order of the collisional transport. Figure 1 illustrates the theoretical prediction of the fluctuation level as a function of the pressure gradient, G_0. Explicit multifold forms of $\hat{\phi}(G_0)$ and $\chi(G_0)$ are seen. The lower-amplitude branch is thermodynamically unstable. Anomalous transport is predicted to occur due to the subcritical excitation, if G_0 exceeds the critical value G_*, which is much smaller than the linear stability boundary G_c.

This analytic estimation is compared to the numerical simulation. Direct nonlinear simulation of the basic set of equations was performed [6]. The two-dimensional turbulence has been calculated in a system of the size $|x| < L_x$ and $|y| < L_y$. (The surface x=0 is the mode rational surface, and z-axis is in the direction of the magnetic field at x = 0. Parameters in the simulation were: $\mu_c = \chi_c = 0.2(c/a\omega_p)^2$, $\mu_{e\,c} = 0.01(c/a\omega_p)^2$, s = 0.5, $L_x = 40(c/a\omega_p)$ and $L_y = 6.4\pi(c/a\omega_p)$.) For this system, the linear stability boundary is given as $G_c \simeq 0.4$, i.e., $\Im_c \simeq 0.25$. Nonlinear excitation of the fluctuations was confirmed in the simulation. Figure 2 compares the theoretical result (the solution of Eqs.(4) and (5), the solid line) of the transport coefficient and that from the numerical simulation (black points). We see that the steady state turbulence is realized even below the critical pressure gradient against the linear

instability, $G_c \simeq 0.4$. The theoretical formula (5) well reproduces the subcritical nature of the turbulence which is obtained by the nonlinear simulation.

4. Summary and Discussion

In summary, the nonlinear theory of the current-diffusive interchange mode turbulence in confined plasmas was developed. A nature of the subcritical turbulence was shown from the theoretical formula. Comparison study with the result from the direct nonlinear simulation was made. The critical pressure gradient for the transition from collisional transport to turbulent transport, $G_0 \geq G_*$, was obtained. The analytic form of the turbulent transport coefficient, Eq.(8), is consistent with the result from the scale invariance technique of the original nonlinear equation [13]. The critical gradient G_* is given in experimental variables as $a/L_p^* = (\varepsilon^2/2\Omega^2)(s^2 g m_i/m_e)^{2/3}(a\nu_i/v_{thi})^{2/3}$, where $1/L_p \equiv -\beta'/\beta$, g is the numerical factor introduced as $\chi_c = g\nu_i\rho_i^2$, and ν_i, ρ_i and v_{thi} are collision frequency, gyro radius and thermal velocity of ions. In the usual experimental circumstances, the condition $G_0 \geq G_*$ is satisfied; The relevance of Eq.(8) with experimental observation is given in [8]. The finding in this article replaces the conventional view of the turbulence theory, in which the forward bifurcation was predicted [3]. The new analytic insight deepens the understanding of the turbulence driven by the pressure gradient in toroidal plasmas, and develops a transport theory for the nonlinear and far-non-equilibrium systems. The pressure gradient coupled to the bad magnetic curvature plays a role of the order parameter to characterize the state of nonuniform plasma. Similar result is obtained for the current diffusive ballooning mode turbulence, which is relevant to tokamaks [11]; the extension will be discussed in elsewhere.

The particular forms of χ and G_* may depend on the choice of the fluid description of the plasma [9], but the basic nature of subcritical excitation is not limited by this assumption. The mechanism for the subcritical excitation is very general, i.e., the parallel electron motion is impeded by fluctuations. (Similar mechanism was also studied in relation to the drift mode [14].) The normal cascade of the spectrum by the

electron nonlinearlity is important. If collisional resistivity alone is taken into account [15], the nonlinear excitation was not obtained. Recently theoretical efforts are also made for the ion temperature gradient (ITG) modes [16]. The impact of the subcritical excitation for such mode will be of interest. Finally, the result of Fig.2 is based on the two-dimensional (2-D) simulation. The 2-D model does not limit the relevance, because the radial mode structure is localized, by the shear, compared to the typical mode separation distance, $1/sk_\theta$; The nonlinear normal cascade by electron dynamics exists in both the 2-D and three-dimensional (3-D) models. Three-D simulations become possible these days [17]. A future test by use of 3-D simulation would be fruitful. These problems are left for the future study.

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Figure Captions

Fig.1 Turbulent transport coefficient (solid line) and fluctuation level (dashed line) as a function of the pressure gradient parameter G_0 which is normalized to $(sa\omega_p/c)^{4/3}\chi_c^{2/3}$. Thin dotted line shows the level of collisional transport. χ_N and $\hat{\phi}$ are normalized to χ_c. (Parameters are $\mu_c = \chi_c = \sqrt{m_i/m_e}\mu_{e,c}$ and $m_i/m_e = 1836$. $\Im_c \simeq 0.25$)

Fig.2 Theoretical prediction for the turbulent transport coefficient (solid line) is compared to the result of the direct nonlinear simulation (dot with error bar). (χ is normalized to $(c/a\omega_p)^2$.) Subcritical appearance of the turbulent transport above the collisional one is confirmed by the simulation. Mixing length estimate, $\chi^{mixing} = \gamma_L/k_\perp^2$, is also shown for comparison (γ_L : linear growth rate).

Fig.1

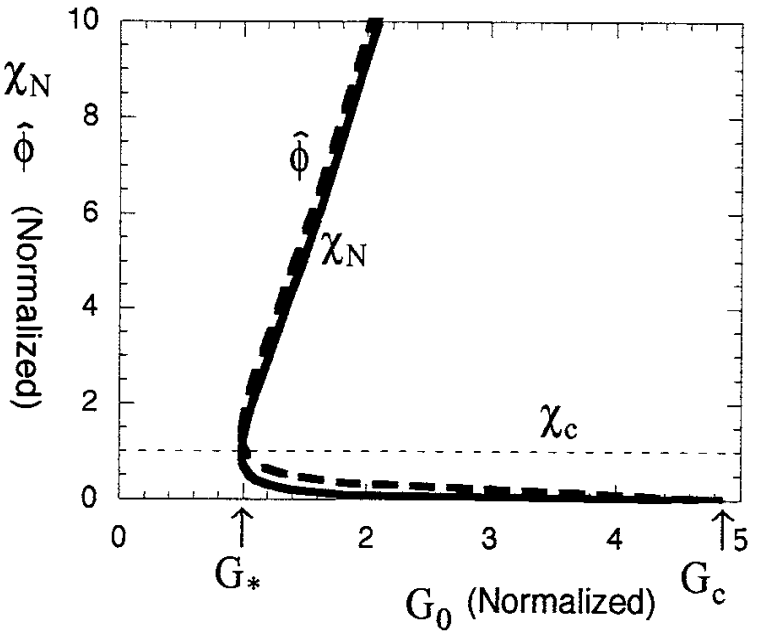
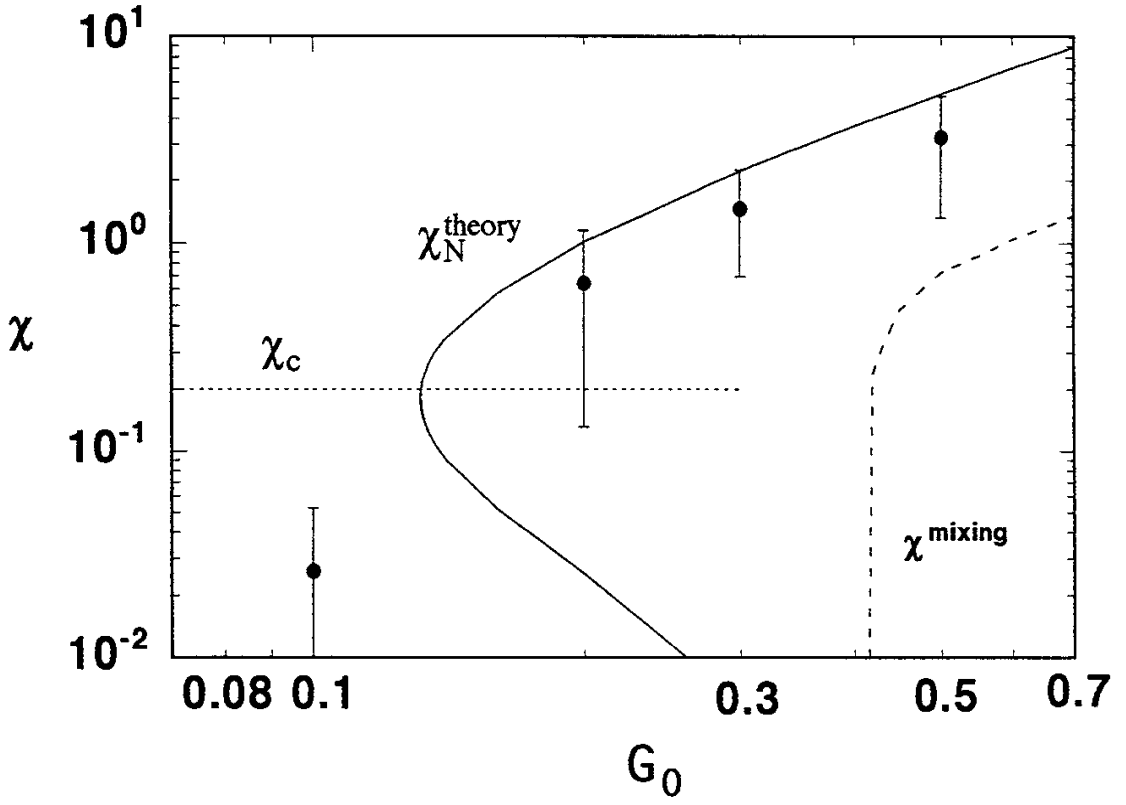

10
8
6
4
2
0
χ_N
$\hat{\phi}$ (Normalized)
$\hat{\phi}$
χ_N
χ_c
0
2
3
4
5
G_*
G_0 (Normalized)
G_c

Fig.2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10^0
10^{-1}
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χ
χ_N^{theory}
χ_c
χ^{mixing}
0.08
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0.5
0.7
G_0

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