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A novel turbulence trigger for neoclassical tearing modes in tokamaks

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Abstract

Stochastic trigger by microturbulence for neoclassical tearing mode (NTM) is studied. NTM induces topological change of magnetic structure and has subcritical nature. Transition rate of, probability density function for and statistically-averaged amplitude of NTM are obtained. Boundary in the phase diagram is determined as the statistical long time average of the transition conditions. NTM can be excited by crossing this boundary even in the absence of other global instabilities.

Keywords: neoclassical tearing mode, statistical theory, stochastic excitation, turbulent noise, probability density function, phase boundary, transition rate, nonlinear instability, long time average

Magnetized plasmas are non-uniform and far from thermal equilibrium. Consequently, various kinds of bifurcations can appear[1,2] producing an abrupt change of the topological structure of the magnetic field. In tokamak and in other laboratory plasmas, such a process appears as a magnetohydrodynamic (MHD) instability named tearing mode.[3,4] It is associated with magnetic field reconnection. Global perturbations with wave numbers perpendicular to the magnetic field are unstable and, due to the plasma resistivity, they can develop radial components that break the field lines. An important problem is now investigated: the possibility of appearing such magnetic surface breaking in ideally stable, low resistivity plasmas.

One possible mechanism is based on a nonlinear instability, the neoclassical tearing mode (NTM).[5-7] This is a subcritically excited tearing mode under the influence of the pressure gradient. The experiments have shown that such perturbations with finite amplitude become unstable even for parameters corresponding to linear stability[8-10] and that they can be triggered by other global MHD instability (as the sawtooth).[10,11] But, in some experiments, the excitation of this instability was produced in the absence of the above trigger.[11,12] The NTM can be stochastically triggered in these conditions. The onset conditions of the NTM are not yet clarified, although the suppression of this instability is necessary for stationary operation of high temperature plasma.[13] The rate of stochastic transition was determined at thermal equilibrium by evaluating the potential barrier crossing induced by thermal fluctuations.[14] It is expected that in nonequilibrium and turbulent plasmas the transition is triggered by the turbulence but there is no theoretical prediction for the excitation rate of the NTM.

In this article, we formulate a Langevin equation for NTM as a stochastic equation in the presence of noise source induced by background fluctuations. The statistical properties of NTM amplitude, such as the probability density function (PDF), the rate of excitation, the average of amplitude, the boundary in the phase diagram and its expression, are derived. We show that the stochastic excitation of NTM is possible to occur without seed island if $\beta_p > \beta_{p*}$ holds. (β_p is the plasma pressure normalized to the poloidal magnetic field pressure.) We note that this mechanism is rather general. For instance, in fluid dynamics it explains the transition of a linearly stable systems in a laminar state (flow in a pipe) to a self-sustained turbulent state.[15] The transition is triggered by random disturbances such as inlet conditions.

The nonlinear instability of the NTM has been discussed, and a dynamical equation of Ohm's law was formulated for the evolution of the amplitude as a deterministic variable

$$\frac{1}{\eta}\frac{\partial}{\partial t}A + \Lambda[A]\,A = 0\,, \qquad (1)$$

where $A \equiv \tilde{A}_* \, q^2 R / B r_s^3 q'$ is the normalized amplitude of the (m, n) -Fourier component of helical vector potential perturbation $\tilde{A}_*$ at the mode rational surface, $r = r_s$, and $-\Lambda$ is the nonlinear growth rate ($-\Lambda > 0$ if unstable). The safety factor $q = rB/B_p R$ as a topological index satisfies the condition $q(r_s) = m/n$ at $r = r_s$. B is the main magnetic field strength, r and R are minor and major radii of torus, $q' = \mathrm{d}q/\mathrm{d}r$, and m and n are poloidal and toroidal mode numbers, respectively. The time is normalized to poloidal Alfven transit time, $\tau_{\mathrm{Ap}} = qR/v_A$ (v_A : Alfven velocity) and the length to r_s . The magnetic island width w , being normalized to r_s , is expressed as $w = A^{1/2}$. The coefficient η is the inverse of resistive diffusion time $\eta = \eta_{//} \mu_0^{-1} r_s^{-2} \tau_{\mathrm{Ap}} = S^{-1}$, where $\eta_{//}$ stands for a parallel resistivity, and S is the Lundquist number.

An explicit form of the growth rate is given by

$$-\Lambda = 2\Delta' A^{-1/2} - \frac{C_1}{W_1^2 + A^2} + \frac{C_2}{W_2 + A} \;, \tag{2}$$

within the neoclassical transport theory, where the first, second and third terms of RHS stand for the effects of current density gradient, polarization drift and bootstrap current, respectively. The term W_1 represents the cut-off due to the banana orbit effect,[16] and we choose a simple model, $W_1 = \rho_b^2 r_s^{-2}$. W_2 represents the cut-off determined by the cross-field energy transport.[17] $C_1 = 2a_{\mathrm{bs}} \beta_{\mathrm{p}} \varepsilon^{1/2} \rho_b^2 r_s^{-2} L_q^2 L_p^{-2}$ and $C_2 = 2a_{\mathrm{bs}} \beta_{\mathrm{p}} \varepsilon^{1/2} L_q L_p^{-1}$ for the limit of small collisions,[7, 18,19] ρ_b is the banana width, L_q and L_p are the gradient scale lengths of safety factor and pressure, respectively, ε is the inverse aspect ratio and a_{bs} is a numerical constant. The parameter Δ' controls the linear stability of tearing mode when induced by the current density gradient.[3, 4] When the amplitude A takes finite values, $-\Lambda$ can be positive even if $\Delta' < 0$, because C_1 and C_2 can be positive. Namely, the mode is nonlinearly unstable while it is linearly stable. Figure 1 illustrates the growth rate as a function of A for the case of $\Delta' < 0$. The marginal stability condition $\Lambda = 0$ can have three solutions at $A \simeq 0$, $A = A_{\mathrm{m}}$ and $A = A_{\mathrm{s}}$ ($A_{\mathrm{m}} < A_{\mathrm{s}}$). A_{m} and A_{s} are the threshold and saturation amplitudes, respectively. Near the linear stability boundary, $\Delta' \simeq 0$, they can be estimated as $A_{\mathrm{m}} = C_1 C_2^{-1}$ and $A_{\mathrm{s}} \simeq C_2^2 / 4\Delta'^2$.

The helical perturbation is subject to a random excitation from the micro turbulent noise. The level of noise is evaluated from the Lagrangian nonlinearity terms, and a stochastic equation is obtained instead of the deterministic equation (1)

$$\frac{\partial}{\partial t} A + \eta \Lambda A = s \frac{\delta^2}{a^2} \left[\phi_h, \Delta A_h \right]_k - s \left[\phi_h, A_h \right]_k - s \frac{v_{\mathrm{Te}}}{v_{\mathrm{A}}} \frac{\delta^2}{a^2} \left[A_h, \Delta A_h \right]_k \;, \tag{3}$$

where $s = aq'/q$ and δ is the collisionless skin depth c/ω_{pe} . ϕ_h is the stream function and A_h is the vector potential of the microscopic turbulence.[20] The suffix h stands for

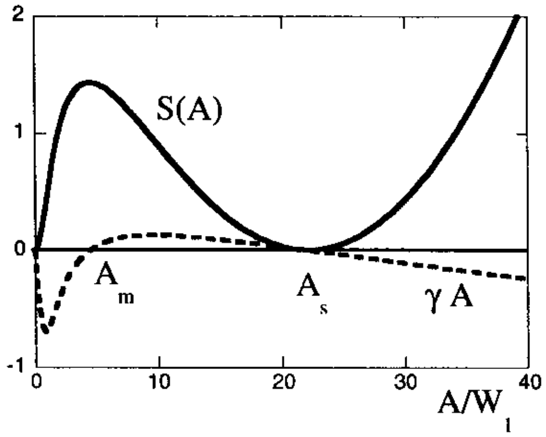

Fig.1 Normalized growth rate multiplied by amplitude, $\gamma A \equiv \Lambda A / C_2$, is shown by dashed line. Zeros indicate the nonlinear marginal stability conditions for the deterministic model. Normalized nonlinear potential $S(A) / \Gamma C_2 W_1$ is shown by the solid line. (Parameters are: $W_1 = W_2$, $C_1/2C_2W_1 = 1$ and $\Delta' W_1^{1/2}/C_2 = -0.0922$.)

the high mode numbers. The Poisson bracket $[u, v]$ is defined as $(\nabla u \times \nabla v)\cdot \boldsymbol{b}$, and $\boldsymbol{b} = \boldsymbol{B}/B$. $[\cdots]_k$ indicates the Fourier component that matches to the test macro mode, and k is the wave number for the macro mode.

We employ the following hierarchical approach. The microscopic turbulence has much shorter autocorrelation time τ_{ac} than that of the global perturbation. They are statistically independent, and the adiabatic approximation is taken. It is induced by plasma pressure gradient being in the nonlinearly marginal state.[21] The saturation levels ϕ_h and A_h can depend on A . We do not consider such dependence here, but it can be introduced in the model.

The RHS of Eq.(3) has two components. One is a coherent part, which has a fixed phase with respect to global perturbation A . The coherent part is renormalized by introducing turbulence-driven transport coefficients such as turbulent resistivity and viscosity. They would modify η , C_1 and C_2 .[19] We note that the sign of C_1 and C_2 can be changed by this renormalization. The electric induction by microfluctuations has been studied in conjunction with dynamo. The $\alpha-$, $\beta-$ or $\gamma-$ dynamo have been known.[1] In this article, however, we use Eq.(2) as a starting assumption and leave the effects of turbulence on Λ for future studies.

The other is an incoherent part. The relative phase to A changes rapidly in time, and contributes to the noise term, being approximated to be random, i.e., $\mathfrak{I}(t) = g\, w(t)$, where g is the magnitude and $w(t)$ indicates white-noise. $\mathfrak{I}(t)$ has a quadratic form of ϕ_h and A_h , and the local instantaneous amplitude of $\mathfrak{I}(t)$ is given as $k k_h^3\, C\, A_h^2$; where numerical constant $C = -s\, f\left(\delta^2 r_s^{-2} + k_h^{-2}\right) + s\sqrt{\beta m_i/m_e}\,\delta^2 r_s^{-2}$ with $f \equiv \phi_h/A_h$ is introduced. Estimations are made as $\Delta A_h = -k_h^2 A_h$ and $|\nabla A_h| = k_h A_h$ for microscopic turbulence, and as $|\nabla A| = k A$ for macro test mode. k_h is the typical mode number of the micro fluctuations, the inverse of which is separated from the coherence length of macro mode. (For a case of ballooning mode turbulence in tokamaks, f is evaluated in ref.21 and is of the order unity.) The statistical average $\sqrt{g^2}$ is related to $|\mathfrak{I}|$ by the law of large numbers. Within the coherent area of global test mode, ℓk^{-1}, a large number ($N = k_h^2 \ell k^{-1}$) of independent kicks contribute to $\mathfrak{I}(t)$. (ℓ : radial scale length of the macro mode. N is evaluated by noting a quasi-two-dimensional feature of fluctuations.) The average $\sqrt{g^2}$ is $N^{-1/2}$ times smaller than the instantaneous local value of $|\mathfrak{I}|$. The magnitude g is evaluated as

$$g^2 = k\, k_h^{-2} \ell^{-1} |\mathfrak{I}|^2\, \tau_{ac} = \ell^{-1} k^3 k_h^4\, C^2 A_h^4\, \tau_{ac}\ , \tag{4}$$

having a dependence like $g^2 \propto \left(\tilde{B}_{r,h}/B_\theta\right)^4 \tau_{ac}$. Experimental magnitude is explained later. The stochastic equation of NTM amplitude A is rewritten as

$$\frac{\partial}{\partial t} A + \eta\, \Lambda\, A = g\, w(t)\ , \tag{5}$$

and A is now a stochastic variable. The statistical property of the NTM amplitude A is studied. It is worthwhile to compare it with Kramers' idea for thermal equilibrium.[14] In Eq.(5), there is a nonlinear force but no Einstein drag term common in Brownian theory; the fluctuations from turbulence are decidedly non thermal unlike standard Langevin theory. The Fokker-Planck equation of $P(A)$ is deduced from Eq. (5) as

$$\frac{\partial}{\partial \tau} P + \frac{\partial}{\partial A}\left(\eta\, \Lambda + \frac{1}{2}\, g\, \frac{\partial}{\partial A}\, g\right) P = 0\ . \tag{6}$$

The stationary solution $P_{eq}(A)$ is expressed as $P_{eq}(A) \propto g^{-1} \exp\left(-S(A)\right)$ by use of a nonlinear dissipation function as $S(A) = \int_0^A 2\eta\, \Lambda(A') g^{-2} A'\, dA'$ which is proportional to the entropy production rate near the thermal equilibrium.[1] Using Eqs.(2) and (4), one has

$$S(A) = \Gamma\left(-\frac{4}{3}\Delta' A^{3/2} + \frac{1}{2} C_1 \ln\left(1 + \frac{A^2}{W_1^2}\right) - C_2\left(A - W_2 \ln\left(1 + \frac{A}{W_2}\right)\right)\right), \tag{7}$$

with $\Gamma = 2 S^{-1} \ell k^{-3} k_h^{-4} C^{-2} A_h^{-4} \tau_{ac}^{-1}$. The coefficient Γ shows a characteristic value of the ratio between the dissipation for crossing over the barrier and excitation by turbulence noise. Its magnitude and dependence are discussed at the end of this article. The PDF is given as $P_{eq}(A) \propto \exp\left(\Gamma \frac{4}{3} \Delta' A^{3/2} + \Gamma C_2 A\right)\left(1 + \frac{A^2}{W_1^2}\right)^{-\Gamma C_1/2}\left(1 + \frac{A}{W_2}\right)^{-\Gamma C_2 W_2}$. The PDF has a stretched non-Gaussian exponential form with power-law dependence. The exponential term is determined by the damping by current density gradient and the drive by bootstrap current. The power-law decay is mainly due to the polarization drift effect. The minimum of $S(A)$, i.e., zero of Λ, predicts the peak of PDF and the probable value of A .

For the case of a bistable state, the nonlinear potential $S(A)$ is shown by solid curve in Fig.1, which has two minima at $A = 0$ and $A = A_s$, separated by a local maximum at $A = A_m$. Statistical transitions take place between these solutions. The dominant (i.e., the most probable) state is determined by the balance between the transition for excitation (from $A = 0$ to $A = A_s$) and the decay (from $A = A_s$ to $A = 0$). The long time average, i.e., the statistical average $\langle A \rangle$, is calculated from the PDF.

Calculating a flux of probability density from Fokker-Planck equation (6),[1, 22] the frequencies of excitation and decay are expressed as

$$r_{ex} = \frac{\eta \sqrt{\Lambda_0 \Lambda_m}}{2\pi} \exp\left(-S(A_m)\right) , \tag{8a}$$

$$r_{dec} = \frac{\eta \sqrt{\Lambda_s \Lambda_m}}{2\pi} \exp\left(S(A_s) - S(A_m)\right), \tag{8b}$$

respectively, where the time rates $\Lambda_{m,s}$ are given as $\Lambda_{m,s} = 2A \left| \partial\Lambda / \partial A \right|$ at $A = A_m$ and $A = A_s$. $P_{eq}(A)$ has a peak at $A = 0$. A noise level where the NTM is not excited is evaluated from a local average of A near $A = 0$, being given as $\langle A_0 \rangle \sim 0.5(-\Gamma\Delta')^{-2/3}$, and yields $\Lambda_0 = \Lambda(\langle A_0 \rangle)$. Note that $\Lambda_{0,m,s}$ are normalized, being of the order unity.

The long time average is given as $\langle A \rangle = \left(A_s r_{ex} + \langle A_0 \rangle r_{dec}\right)\left(r_{ex} + r_{dec}\right)^{-1}$. $\langle A \rangle$ approaches to A_s if $r_{ex} > r_{dec}$ holds. It reduces to $\langle A_0 \rangle$, if $r_{ex} < r_{dec}$ holds. The phase boundary for the statistical average is determined by the condition $r_{ex} = r_{dec}$. Apart from a logarithmic dependence, the condition is given by $S(A_s) = 0$. Figure 2 shows the statistical average $\langle A \rangle$, together with threshold and saturation amplitudes (A_m and A_s), as a function of β_p . $\langle A \rangle$ drastically changes across the condition $\beta_p = \beta_{p*}$, a formula of which is derived as follows. From Eq.(7), the condition $S(A_s) = 0$ is rewritten as $-\frac{4}{3} \Delta' A_s^{3/2} = C_2 A_s - \frac{1}{2} C_1 \ln\left(A_s^2 W_1^{-2}\right)$ where $A_s >> W_1, W_2$ is assumed. Using the relation $A_s \simeq C_2^2 / 4\Delta'^2$, we have

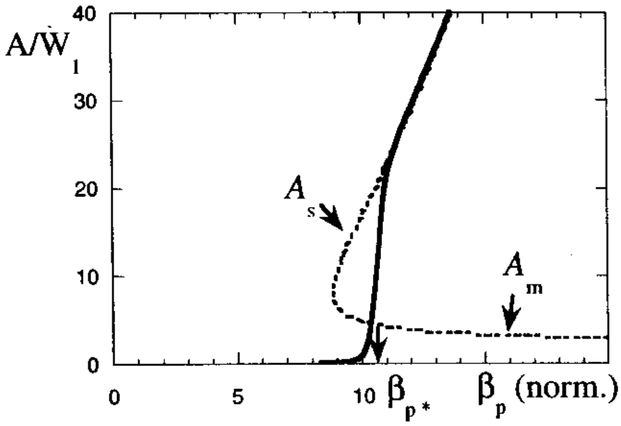

Fig.2 Amplitude of NTM as a function of the plasma pressure. Solid line shows the statistical average $\langle A \rangle$. A thin dotted line indicates the threshold $A_{\rm m}$ and saturation amplitude A_s for the deterministic model. Normalized $\beta_{\rm p}$ is $C_2 / \left(-\Delta' W_1^{1/2} \right)$, i.e., $\left(2 a_{\rm bs} \varepsilon^{1/2} L_q r_s / \rho_b L_p (-\Delta') \right) \beta_{\rm p}$. (Parameters are: $W_1 = W_2$, $C_1 / 2 C_2 W_1 = 1$, $\Gamma C_2 W_1 = 5$.)

$$\Delta' = \Delta'_* \equiv -\sqrt{\frac{1}{12}}\, C_2^{3/2} C_1^{-1/2} \left(\ln \left(3 C_1 / 2 C_2 W_1 \right) \right)^{-1/2} , \tag{9}$$

where dominant terms are retained. The boundary Δ'_* is negative and of the order unity. Equation (9) is rewritten as $\Delta' = -\sqrt{\varepsilon L_q / 3 L_p} \left(\ln \left(3 L_q / 2 L_p \right) \right)^{-1/2} a_{\rm bs}\, r_s \rho_b^{-1} \beta_{\rm p}$ by substituting C_1 and C_2. It is reformulated in a form of a critical pressure as $\beta_{\rm p} = \beta_{\rm p*} = \sqrt{3 L_p / \varepsilon L_q} \left(\ln \left(3 L_q / 2 L_p \right) \right)^{1/2} a_{\rm bs}^{-1} (-\Delta')\, \rho_b r_s^{-1}$.

An example of the transition frequency is estimated in the following. Near the linear stability condition, $\Delta' \simeq 0$, one has $S(A) = \left(C_1 W_1^{-2} - C_2 W_2^{-1} \right) A^2/2 + C_2 W_2^{-2} A^3/3 - \left(C_1 W_1^{-4} + C_2 W_2^{-3} \right) A^4/4 + \cdots$. The potential barrier $S(A_m)$ is given by the maximum. For the case of $W_2 > W_1$, one has a simple estimate $S(A_m) \simeq \left(1 - 2 C_2 W_1^2 / C_1 W_2 \right) C_1/4 \sim C_1/4$, by keeping the first order correction of W_1/W_2. Substituting it into Eq.(8a), one gets the excitation rate of NTM as

$$r_{\rm ex} \simeq \frac{\eta \sqrt{\Lambda_0 \Lambda_{\rm m}}}{2\pi} \exp\left(-\Gamma \frac{C_1}{4} \right) . \tag{10}$$

The parameter Γ is the key for the transition frequency. For L-mode plasmas, when one employs the current-diffusive ballooning mode as micro mode, one has

$A_h \simeq 10\, s\, \alpha^2 (\delta/r_s)^2$, $\phi_h \simeq 10\, \alpha^{3/2} (\delta/r_s)^2$ and $\tau_{ac} \sim \alpha^{-1/2}$, where $\alpha = -q^2 R \mathrm{d}\beta/\mathrm{d}r$ is the normalized pressure gradient.[23] Substituting them into the formula of Γ below Eq.(7), one has $\Gamma = 2\ell k^{-3} \left(-\alpha^{-1/2}(1+\alpha) + s\sqrt{\beta m_i / m_e} \right)^{-2} 10^{-4} s^{-4} \alpha^{-11/2} S^{-1} (\delta/r_s)^{-8}$. The argument $\Gamma C_1/4$ in Eq.(10) may be simplified as $4^{-1} a_{bs} \varepsilon^{1/2} L_q^2 L_p^{-2} s^{-2} (m_e/\beta m_i)\, \ell k^{-3} 10^{-4} \alpha^{-11/2} \beta_p S^{-1} \rho_b^2 r_s^6 \delta^{-8}$ for $\beta m_i/m_e > 1$. This result shows that when the resistivity becomes so low as to satisfy the condition $S \simeq 10^{-4} (m_e/\beta m_i)\, \ell\, k^{-3} \alpha^{-11/2} \rho_b^2 r_s^6 \delta^{-8}$, the exponential term becomes of the order of unity, and the transition frequency of the order of η is expected. When the plasma pressure gradient becomes large, a strong turbulence (M-mode) has been predicted.[21, 24] In this case, $\tilde{A}_h$ is enhanced by the factor of $(\alpha \beta_i m_i/m_e)^{1/2}$. One has $\Gamma C_1/4 \simeq 4^{-1} a_{bs} \varepsilon^{1/2} L_q^2 L_p^{-2} s^{-2} (m_e/\beta m_i)^3\, \ell k^{-3} 10^{-4} \alpha^{-15/2} \beta_p S^{-1} \rho_b^2 r_s^6 \delta^{-8}$. The condition of frequent transitions, $\Gamma C_1/4 \sim 1$, is given as $S \simeq 10^{-4} (m_e/\beta m_i)^3\, \ell\, k^{-3} \alpha^{-15/2} \rho_b^2 r_s^6 \delta^{-8}$. This condition might be easily satisfied in a high temperature experiment of modern tokamaks.

In summary, we have developed a statistical theory for the excitation of nonlinear NTM. The stochastic equation is formulated including the subcritical excitation mechanism of NTM. The rate of transition and statistical average of amplitude are derived, and the phase boundary in plasma parameter space, β_{p*} or Δ'_* , is obtained. Linearly stable systems are prone to nonlinear instability if $S(A_s) < 0$ holds. The formula is applied to either cases of micro fluctuations or of other random MHD activities. Experimental database for the presence of NTM must be compared with the result of phase boundary derived from the statistical theory. The rate of stochastic transition depends on the microfluctuation level and is evaluated for example cases. However, the boundary is given by $S(A_s) = 0$ and is insensitive to the magnitude of micro fluctuations. It is plausible that the stochastic transition without the trigger by large MHD events (e.g., sawtooth or fish-bone instabilities) can be observed in high temperature tokamak plasmas if the condition $\beta_p > \beta_{p*}$ is satisfied. This explains observations in refs.11 and 12.

Equation (8) is a generalization of the result of thermal equilibrium, i.e., Eq.(476) of ref.14 that recovers Arrhenius' law, to the case of the turbulence trigger. The turbulence amplitude is included in the denominator of $S(A)$ that appears in exponential term of r_{ex} and r_{dec} . Owing to the turbulence trigger, the transition probability is greatly enhanced and the variation of the average $\langle A \rangle$ across $\beta_p = \beta_{p*}$ becomes less sharp. The energy of microfluctuations is estimated in tokamak turbulence and is about $\delta^2 r_s \lambda_D^{-3}$ times larger than that in thermal equilibrium (§23 of ref.1). In the latter case, Γ is larger by a factor $\delta^4 r_s^2 \lambda_D^{-6}$ and the transition is very difficult to occur.

This article does not give complete picture for the trigger of NTM but provides a theoretical framework for future studies. There are a lot of effects and contributions which could be incorporated in the nonlinear statistical theory. (Examples include: The coherent part of RHS of Eq.(3), like dynamo term and other nonlinear drags, can influence Λ so as to modify β_{p*} ; Excitation of large scale island, in turn, may suppress the transport as in the case of Snakes.[25]; Semi-micro structures could coexist as reviewed in ref.1.) The analytic formula (9) and (10) could be verified by direct solution of Eq.(5) by Monte Carlo simulation. These are left for future studies and will give quantitative results.

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