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Edge and divertor plasma: detachment, stability, and plasma-wall interactions

S.I. Krasheninnikov¹, A.S. Kukushkin^{2,3}, Wonjae Lee¹, A.A. Phsenov^{2,3}, R.D. Smirnov¹, A.I. Smolyakov⁴, A.A. Stepanenko³ and Yanzeng Zhang¹

¹ University of California San Diego, La Jolla, CA 92093, United States of America

² Kurchatov Institute, Kurchatov sq. 1, 123182 Moscow, Russian Federation

³ NRNU MEPhI, Kashirskoje av. 31, 115409 Moscow, Russian Federation

⁴ University of Saskatchewan, 116 Science Place, Saskatoon, SK S7N 5E2, Canada

E-mail: skrash@mae.ucsd.edu

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Abstract

The paper presents an overview of the results of studies on a wide range of the edge plasma related issues. The rollover of the plasma flux to the target during progressing detachment process is shown to be caused by the increase of the impurity radiation loss and volumetric plasma recombination, whereas the ion-neutral friction, although important for establishing the necessary edge plasma conditions, does not contribute *per se* to the rollover of the plasma flux to the target. The processes limiting the power loss by impurity radiation are discussed and a simple estimate of this limit is obtained. Different mechanisms of meso-scale thermal instabilities driven by impurity radiation and resulting in self-sustained oscillations in the edge plasma are identified. An impact of sheared magnetic field on the dynamics of the blobs and ELM filaments playing an important role in the edge and SOL plasma transport is discussed. Trapping of He, which is an intrinsic impurity for the fusion plasmas, in the plasma-facing tungsten material is considered. A newly developed model, accounting for the generation of additional He traps caused by He bubble growth, fits all the available experimental data on the layer of nano-bubbles observed in W under irradiation by low energy He plasma.

Keywords: edge plasma, divertor detachment, current convective instability, blobs, He nano-bubbles

(Some figures may appear in colour only in the online journal)

1. Introduction

The processes involving the edge plasma and plasma-material interactions in magnetic fusion devices are very multifaceted and include a wide spectrum of phenomena ranging from plasma turbulence and meso-scale stability, impurity and radiation transport, recycling and transport processes of hydrogen species in the wall material, to modification of the wall material properties and the surface morphology. In many cases, these processes are strongly coupled and exhibit synergistic effects, which make it difficult to analyze and interpret the edge plasma phenomena observed in experiments.

In this overview paper we present the results of our studies on a wide range of issues related to the edge plasma. In section 2 we report on the results of our numerical simulations closing the long-standing dispute on the roles of impurity radiation, volumetric plasma recombination and ion-neutral friction in the rollover of the plasma flux to the target, which is the manifestation of detachment. We show that the rollover is caused by the increase of the impurity radiation loss and volumetric plasma recombination whereas the ion-neutral friction, although important for establishing the necessary edge plasma conditions, does not contribute *per se* to the rollover of the plasma flux to the target. We find that expansion of the region of the cold, recombining plasma in divertor by detachment does naturally limit the amount of power that can be radiated by impurities. In section 3 we identify different mechanisms of meso-scale thermal instabilities driven by impurity radiation and resulting in self-sustained oscillations in the edge plasma. These oscillations cause a significant ~30-50% variation of the heat load on the divertor targets. In section 4 we show that the current-convective instability can develop in detached inner divertor whereas the outer one is still attached. This can explain the fluctuations of impurity radiation in the detached inner divertor observed in the AUG tokamak. In section 5 we demonstrate that the plasma blobs seen in many experiments both inside and outside the separatrix can be described within the framework of nonlinear drift wave physics. Section 6 is devoted to the study of blob dynamics in a sheared magnetic field. We show, in particular, that the difference in the plasma density at the sheaths results in a formation of the electrostatic potential that can cause motion of the filaments even in the absence of any source of plasma polarization. Finally, in section 7 we consider the dynamics of formation of the He nano-clusters in the plasma-facing tungsten material. Our newly developed model accounting for generation of additional He traps caused by He bubble growth fits all the available experimental data on the layer of nano-bubbles observed in W under irradiation by low energy He plasma.

2. Roles of plasma-neutral interactions, impurity radiation, and plasma recombination in the divertor plasma detachment

Although the detached divertor regime was discovered more than two decades ago (see [1] and the references therein) and is now considered the primary operation regime for ITER and future tokamak reactors, the cause of the rollover of a total plasma flux to the divertor targets, Γ_r , (see e.g. figure 2 in [1]), which is the signature of the detachment phenomenon, is still under debate. Resolution of this issue requires clear understanding of the physics underlying the plasma detachment phenomena.

Two different explanations of the rollover have been put forward. In [2], the reduction of the plasma flux to the targets was attributed to ion-neutral 'friction', which at low divertor plasma temperature switches the plasma flow regime in the recycling region from free streaming to diffusive-like. However, this approach has many uncertainties related to the difference between plasma and neutral velocities (which is a dense plasma can be small) and momentum transport to material surfaces. Therefore, in [3, 4] the rollover was analyzed from the point of view of particle and energy balances, which have a solid background. For relatively low plasma temperature in the vicinity of the targets, T_t , the results of this analysis can be boiled down to the following expression for Γ_t (see p 10 in [7] for details):

$$\Gamma_{\rm t} = (Q_{\rm SOL} - Q_{\rm imp})/E_{\rm ion}^{\rm eff} - \Gamma_{\rm rec} \equiv \Gamma_{\rm ion}^{\rm satur} - \Gamma_{\rm rec}, \qquad (1)$$

where Q_{SOL} and Q_{imp} are the energy flux from the core into the scrape-off layer (SOL) and the energy loss related to the impurity radiation, E_{ion}^{eff} is the effective neutral 'ionization cost' (including the energy loss associated with hydrogen ionization and elastic collisions per an ionization event), which determine the saturation of the ionization source, Γ_{ion}^{satur} , and $\Gamma_{\rm rec}$ is the plasma recombination sink. Based on equation(1), it was concluded that although the ion-neutral interactions are important for establishing the proper plasma conditions, they are unable to cause any significant reduction of the plasma flux to the target and the actual reasons for the rollover are the impurity radiation loss and the recombination processes. Moreover, only recombination processes can reduce plasma flux to the target below the saturation level, $\Gamma_{\rm ion}^{\rm satur}$, which we will call ultimate detachment.

The conclusions following from equation (1) are consistent with the experimental data (see e.g. [5]). The recombination sink can reach ~ 80% (and even higher) of the ionization source and, in accordance with equation (1), Γ_t can be increased with addition of auxiliary heating or reduced by impurity puffing (see figure 6 from [5]). Here we compare the results of [3, 4] with those from 2D edge plasma modeling to prove the physical correctness of the simplifications made.

We use the SOLPS4.3 code package [6] to model divertor plasma detachment in DIII-D-like plasma (similar to DIII-D magnetic configuration and wall geometry). In our simulations we use a 'closed box' approximation (no puffing nor pumping is applied), where the edge plasma parameters are completely determined by cross-field transport, Q_{SOL} , and fixed content (neutrals and ions) of deuterium, $N_{\rm D}^{\rm edge}$, and impurity, $N_{\rm imp}^{\rm edge}$, species in the simulation domain. An increase in either $N_{\rm D}^{\rm edge}$ or $N_{\rm imp}^{\rm edge}$ with fixed $Q_{\rm SOL}$ causes progressive cooling of the divertor plasma, which leads eventually to ultimate detachment and the rollover of Γ_t . We choose the cross-field particle and heat diffusivities to be constant ($D_{\perp} = 0.3 \text{ m}^2 \text{ s}^{-1}$, $\kappa_{\perp}^{e,i} = 1 \text{ m}^2 \text{ s}^{-1}$). We account for plasma-neutral elastic interactions, neutral excitation and ionization, and plasma recombination effects (more details can be found in [7]). We start with the pure deuterium plasma and perform a scan over $N_{\rm D}^{\rm edge}$ for Q_{SOL} equal to 4 and 8 MW for the cases with recombination turned on or off.

The results of the simulations showing the dependences of $\Gamma_{\rm t}$ on $N_{\rm D}^{\rm edge}$ are presented in figures 1(*a*) and (*b*) and demonstrate a good agreement with the conclusions following from equation (1). As we see in figure 1(*a*), only the case with plasma recombination turned on exhibits a strong decrease of $\Gamma_{\rm t}$ similar to that observed in experiments, while with recombination turned off, $\Gamma_{\rm t}$ virtually saturates with increasing $N_{\rm D}^{\rm edge}$ at the level $\Gamma_{\rm ion}^{\rm satur}$. We notice that for $Q_{\rm SOL} = 4$ MW, $Q_{\rm rad} = 0$, and $E_{\rm ion}^{\rm eff} \sim 30$ eV, equation (1) gives $\Gamma_{\rm ion}^{\rm satur} \sim 8 \cdot 10^{23} \, {\rm s}^{-1}$, which is very close to the results of numerical simulation shown in figure 1(*a*). Adding fixed amount of impurity radiation loss ($Q_{\rm rad} = 4$ MW) for the case with $Q_{\rm SOL} = 8$ MW and no recombination reduces $\Gamma_{\rm ion}^{\rm satur}$ to the level corresponding to the case $Q_{\rm SOL} = 4$ MW and again, only volumetric recombination can cause the ultimate detachment and the rollover of $\Gamma_{\rm t}$ (see figure 1(*b*)).

Thus we see that the rollover of Γ_t is indeed can be caused by impurity radiation and plasma recombination, while ion-neutral interactions *per se* are unable to reduce Γ_t . Nonetheless, the ion-neutral collisions play two important



Figure 1. Plasma flux to the targets for different Q_{SOL} with and without impurities and recombination processes versus $N_{\text{D}}^{\text{edge}}((a) \text{ and } (b))$ and the onset of ultimate detachment (dashed line) versus $P_{\text{up}}/q_{\text{recycl}}(c)$ found from numerical simulations.

roles in detachment physics. First, at low plasma temperature near the targets, the ion-neutral friction/viscosity provides an effective force that counter-balances the upstream plasma pressure. Secondly, when the plasma temperature in the neutral ionization region becomes low (~few eV), neither atomic hydrogen nor impurity-related atomic processes associated with radiation effects can cool the plasma further down to ~eV range necessary to activate plasma recombination (e.g. see strong reduction of impurity cooling function at low temperatures in [47]). Then, the electron-ion-neutral energy exchange (including elastic ion-neutral collisions and vibrational excitation of hydrogen molecules by electron impact, see corresponding rate constants in [48]) assisted by fast heat conduction by neutrals dumping the residual plasma energy to the target becomes important for cooling the plasma down to ~eV range and boosting the recombination processes.

Apart from verification of the conclusions made in [3, 4] with respect to the mechanisms of detachment, we also check the predictions for the onset of an ultimate detachment on a particular flux tube. In [7] it was shown analytically that the 'rollover' of the plasma flux within a specific flux tube, j_i , caused by plasma recombination should occur when the ratio of the 'upstream' plasma pressure, P_{up} , to the specific power flux entering the recycling region along the magnetic field lines, q_{recycl} , exceeds some critical value:

$$P_{\rm up}/q_{\rm recycl} \approx (P_{\rm up}/q_{\rm recycl})_{\rm crit} \sim 15 \,\mathrm{N}\,\mathrm{MW}^{-1},$$
 (2)

here we assume deuterium plasma. Since for a low T_t , impurity radiation largely occurs at the temperatures higher than plasma temperature in recycling region we have $q_{\rm recycl} = q_0 - q_{\rm rad}$, where q_0 is the heat flux from entering divertor and q_{rad} is the impurity radiation loss. The dependences of j_t (for the flux tube adjacent to the separatrix in the inner divertor) on $P_{\rm up}/q_{\rm recycl}$, following from our simulations for different Q_{SOL} and Q_{rad} , that are shown in figure 1(c), confirm this criterion both qualitatively and, with a reasonable accuracy, even quantitatively. Earlier, in [8], it was shown analytically and verified with 2D numerical simulations that once ultimate detachment proceeds, P_{up} and, correspondingly, n_{up} saturate at the level close to that determined by equation (2) with increasing the total number of hydrogen neutral particles and ions in the simulation domain $N_{\rm D}^{\rm edge}$. Simultaneously, the region occupied by cold recombining plasma containing both impurity and neutrals expands and accumulates more and more particles. Moreover, using the expressions from [49] for both drag and thermal forces, we find that plasma flow even with pretty small Mach numbers, exceeding ~ $2.2 \times \lambda_C/L_T \ll 1$ (where λ_C and L_T are the Coulomb mean free path and temperature variation scale length), is already enough to drag impurity toward the targets and confine it in a cold divertor region. The physical reason for such a saturation caused by plasma recombination can be explained as follows. Any significant increase in P_{up} above the value given by (2) will cause an increase in the plasma density in the cold region, which would boost plasma recombination and, as a result, the plasma flow from upstream, thus reducing P_{up} .

However, the impurity radiation loss $q_{\rm rad}$, changing $q_{\rm recycl}$, depends on both the impurity and plasma densities, while the latter is limited by equation (2). Thus, equation (2) is actually a nonlinear equation limiting both the upstream plasma density $n_{\rm up} \approx n_{\rm up}^{\rm max}$ and the impurity radiation loss $Q_{\rm imp} \approx Q_{\rm imp}^{\rm max}$. Theoretical assessment of the dependencies of $n_{\rm up}^{\rm max}$ and $Q_{\rm imp}^{\rm max}$ on $Q_{\rm SOL}$, plasma current, $I_{\rm p}$, toroidal magnetic field, $B_{\rm T}$, minor and major tokamak radii, *a* and *R*, includes models for both the impurity radiation loss and anomalous plasma transport. In the simplest form such a model includes the energy balance equation for the SOL, which, assuming that the main energy transport channel is heat conduction, can be written as

$$n_{\rm up}\kappa_{\perp}T_{\rm up}/\Delta_{\rm SOL}^2 = \hat{\kappa}_{\parallel}T_{\rm up}^{7/2}/\ell_0^2 \equiv q_0/\ell_0,$$

$$Q_{\rm SOL}/S_{\rm sep} = n_{\rm up}\kappa_{\perp}T_{\rm up}/\Delta_{\rm SOL},$$
(3)

where κ_{\perp} and $(2/7)\hat{\kappa}_{\parallel}T_{up}^{5/2}$ are the anomalous perpendicular and classical parallel heat diffusivity and heat conduction coefficient, Δ_{SOL} is the characteristic width of the SOL, ℓ_0 the parallel connection length and S_{sep} the area of the last closed flux surface. To describe the anomalous cross-field transport, we adopt the 'heuristic' model from [9], which can be written as $\Delta_{SOL} = R(Q_{SOL}/\hat{Q})^{1/8}$, where $\hat{Q} \propto I_p^9 R^{16}/a^{17} B_T^2$ is the normalization parameter. For the radiation loss we assume $q_{rad} = \xi_{imp} L_{imp} n_{up}^2 \ell_0$, where ξ_{imp} and L_{imp} are the impurity fraction and effective cooling function, which we assume to be constant. Then from equation (3) we find $Q_{rad}^{max}/Q_{SOL} = F(\Phi)$, where $F(\Phi) \approx 1 - 1/\Phi$ for $\Phi > 1$, $F(\Phi) \approx \Phi^2$ for $\Phi < 1$, and

$$\Phi \propto \xi_{\rm imp}^{1/2} ({\rm RI}_{\rm p}/Q_{\rm SOL})^{3/16} B_{\rm T}^{15/64} (R/a)^{27/112}.$$
 (4)

As follows from equation (4), Q_{rad}^{max}/Q_{SOL} shows a rather weak dependence on all input parameters. However, more work is needed to verify these conclusions.



Figure 2. Maximum electron temperature at the target and impurity radiation loss versus Ne content (*a*) and the target peak heat load variation caused by the SSOs driven by high-Z(b) and low-Z impurity radiation (*c*) [14]. Reproduced from [14], with the permission of AIP Publishing.

3. Stability of detached plasma and meso-scale thermal instabilities

There is a significant body of analytic and semi-analytic studies devoted to the stability of highly radiative divertor plasmas (see e.g. [10–12]), which indicate that the transition to the detached divertor regime may be bifurcation-like. However, the models used in these studies employ a simplified description of both the plasma and impurity dynamics (e.g. a fixed impurity fraction), which in practice may be not valid. Therefore, we perform 2D numerical simulations addressing this issue for the DIII-D-like plasmas (see [13] for details). We start with fixed cross-field transport coefficients (the same we used in previous Section) and $N_{\rm D}^{\rm edge} = 3.5 \times 10^{20}$ for $Q_{\text{SOL}} = 8 \,\text{MW}$ (which, as follows from figure 1, is below the rollover of Γ_t for the pure deuterium plasma) and then gradually increase the neon impurity content, $N_{\rm Ne}^{\rm edge}$, in our simulation domain. The results show no trace of bifurcation-like transitions and in figure 2(a) one can see a smooth increase in $Q_{\rm rad}$ and decrease in the maximum electron temperature in the outer divertor, $T_{\rm e}^{\rm max}$ along with increasing $N_{\rm Ne}^{\rm edge}$.

The reason for the absence of traces of thermal instability caused by the impurity radiation loss is the following. Although the amount of impurity in the computational domain increases when detachment proceeds, a low temperature region starts to develop in the divertor volume, leaving the upstream plasma parameters at some saturated level (see figure 4 from [8]). This cold plasma does practically not radiate, but a large amount of both deuterium and impurity are residing there, which effectively reduces the amount of actually radiating impurity. Therefore, this cold region works as a stabilizing reservoir: an increase in the radiation loss increases the volume of the cold region that stores the excess of impurity and prevents further plasma cooling.

However, if we take into account that the cross-field plasma transport coefficients can increase when detachment proceeds (which indeed was observed in experiments [15]), then our 2D SOLPS simulations show a bifurcation of the divertor plasma parameters. The reason for this is a positive feedback of widening of the SOL with detachment, which is accompanied by the reduction of q_{recycl} , which further promotes detachment and widening of the SOL [13]. The bifurcation of the plasma parameters can also be related to the breakdown of quasi-equilibrium between the absorption and desorption processes on the plasma facing components (PFC) with the reduction of plasma temperature and, therefore, the energy of ions impinging on the PFCs (a large amount of hydrogen is usually stored in the PFCs material) [13].

Our simulations of semi-detached ITER-like plasma demonstrated development of Self-Sustained Oscillations (SSO) in the edge plasma, driven by low- or high-Z impurity radiation [14]. The high-Z impurity (tungsten)-driven SSO were observed in a proximity to the X-point. Tungsten appeared in the edge plasma because of ablation of the dust particles injected from the divertor targets to simulate a possible impact of ELMs with DUSTT code [16]. The physics of these SSO is related to the effect of the thermal force acting on the impurity ions by radiation-condensation instability (RCI). The thermal force is pushing impurity towards the high temperature regions, which switches the RCI from aperiodic instability to a slowly growing traveling wave [17], which in the nonlinear phase transforms into the SSO. The physics of the SSO driven by low-Z impurity (nitrogen) is somewhat similar to that observed in [18] for the pure hydrogen plasma and can be described as follows. Nitrogen neutrals in the divertor volume are ionized predominantly in a relatively dense and hot plasma region and then move along the magnetic field lines. As a result, both the fraction of the N ions and the radiation loss within that flux tube increase. This causes the reduction of the plasma temperature and N ionization rate within the flux tube. The correlation of the waveforms of plasma cooling, N transport and ionization results in development of the SSOs.

The both high- and low-Z impurity-induced divertor plasma oscillations cause large, \sim 30–50%, variation of the peak heat load on the divertor targets with a period of \sim 10 ms and longer, see figures 2(*b*) and (*c*). This variation can lead to significant periodic changes of the surface temperature and the under-surface temperature gradient in the target material, exacerbating material damage due to the thermal fatigue. Note that the both oscillation mechanisms described rely on the plasma transport phenomena different from MHD instabilities and can present a potential challenge for ITER divertor operation even in the absence of ELMs. However, more studies should be done to identify the conditions at which these oscillations develop.

4. Current-convective instability in detached plasmas

The asymmetry of the inner and outer divertors, which forces the inner divertor to detach first while the outer one is still attached, results in the large temperature difference between the vicinities of the inner and outer targets. This temperature difference causes the onset of a large electric potential drop through the detached plasma of the inner divertor [19]. Large potential drop along with inhomogeneity of the resistivity of the detached plasma across the divertor leg can drive the current-convective instability (see e.g. [20]) in the inner divertor and subsequent fluctuations of the radiation loss similar to that observed in experiments. Usually, the current-convective (or 'rippling') mode is stabilized in tokamaks due to the high electron thermal conductivity along the magnetic field, which quickly smooths out the electron temperature inhomogeneity driving this instability (see e.g. [21]).

However, for the case of the cold (~1 eV) detached plasma, the electron thermal conductivity becomes low and unable to stabilize the current-convective instability completely [22]. Thorough analysis of the current-convective instability for the case of detached inner and attached outer diverters shows that its development is limited by parallel electron heat conduction and magnetic sheath effects, which gives the maximum growth rate ~ $10^4 \div 10^5 \,\mathrm{s}^{-1}$ [22]. This instability causes fluctuations of the plasma pressure in detached plasma along the magnetic field lines, which results in bursts of the plasma flow from radiation region beyond the detachment front down to the divertor targets and, therefore, subsequent fluctuations of radiation losses similar to that observed in experiments [23]. Assuming that in a nonlinear regime the characteristic frequency of the plasma parameter fluctuations is of the order of the growth rate of the current-convective instability, we find a reasonable agreement with experimental data showing ~ 10 kHz frequency of the radiation loss oscillations. Once the outer divertor detaches also, the temperature difference between the vicinities of the inner and outer targets disappears and the driving force for the currentconvective instability, causing oscillation of the radiation losses, vanishes. This feature is indeed observed in experiments [23].

5. On the blob formation mechanism

The intermittent blobby cross-field plasma transport is known to play a crucial role in the transport of the edge plasma in magnetic fusion devices (see e.g. [24, 25]). While the mechanism propelling high plasma density filaments (blobs) on the outer side of a tokamak in the (SOL) is rather well understood [24, 26], the physics responsible for the formation of the blobs is still not clear. However, there is a significant body of experimental data showing that high-density blobs exist already inside the separatrix (see e.g. [27, 28, 50]) where they move mainly in poloidal direction and from time to time cross the separatrix and appear in the SOL. This suggests that the blob formation mechanism is not related to either separatrix or SOL specific features. Recently, by retaining nonlinear electron Boltzmann factor in Hasegawa-Mima equation, it was shown [29, 51] that inherence of formation of the high plasma density blobs to the drift wave dynamics is plausible.

Following [30] we consider the drift wave dynamics assumingBoltzmannelectrons $n_e(\vec{r}, t) = n_0 \exp\{\phi(\vec{r}, t) - \Lambda x\}$, where both n_0 and Λ are constant, $\phi = e\varphi/T_e$, *e* is the electron charge, φ is the electrostatic potential, $T_e = \text{const.}$ is the electron temperature, and *x* and *y* are, respectively, radial and poloidal coordinates. We notice that unperturbed plasma density is proportional to $exp(-\Lambda x)$, so that Λ plays the role of effective inverse radial scale length of unperturbed plasma density. In a simplest liner theory of the drift waves it results in the standard dispersion equation $\omega = (cT_e/eB)\Lambda k_y/(1 + k^2\rho_s^2)$, where ω and \vec{k} are the drift wave frequency and wave vector, and $\rho_s^2 = T_e M(c/eB)^2$. Then, adopting a cold ion approximation, but keeping a nonlinear Boltzmann term, from the quasi-neutrality condition we find a modified Hasegawa– Mima equation with nonlinear electrons:

$$\mathbf{d}_{\mathbf{t}}\mathbf{e}^{\phi} - \rho_{\mathbf{s}}^{2}\nabla\cdot(\mathbf{e}^{\phi}\mathbf{d}_{\mathbf{t}}\nabla\phi) - \Lambda\mathbf{e}^{\phi}\vec{e}_{x}\cdot(\vec{V}_{0} - \rho_{\mathbf{s}}^{2}\mathbf{d}_{\mathbf{t}}\nabla\phi) = 0, \quad (5)$$

where $\vec{V}_0 = -D_B(\nabla \phi \times \vec{b})$, $d_t(...) = \partial_t(...) + (\vec{V}_0 \cdot \nabla)(...)$, $D_B = cT_e/eB$, *M* is the ion mass and *B* is the strength of the magnetic field. One can show that equation (5) conserves two integrals:

$$I_1 = \langle exp(\phi - \Lambda x) \rangle_{x,y} \text{ and } I_2 = \left\langle \phi - \rho_s^2 (\nabla \phi)^2 / 2 \right\rangle_{xy}, \tag{6}$$

where $\langle ... \rangle_{xy}$ denotes averaging over the coordinates x and y.

In what follows, we will show that equation (5) allows the solutions with large amplitude solitary-like plasma density structures resembling blobs. Considering equation (5) for the case where $|\partial(...)/\partial y| \gg |\partial(...)/\partial x|$, that is, looking for streamer-like structures, from equation (5) we find

$$\partial_{t}\mathbf{e}^{\phi} - \rho_{s}^{2}\partial_{y}(\mathbf{e}^{\phi}\partial_{yt}^{2}\phi) + U_{\mathrm{DW}}\partial_{y}\mathbf{e}^{\phi} = 0, \tag{7}$$

where $U_{\text{DW}} \equiv \Lambda D_{\text{B}}$ is the phase velocity of the linear drift waves in the long wavelength approximation. For the case of equation (7) the conservation of I_{I} implies the conservation of the plasma density, averaged over y, at every x. Therefore, we can take $I_{\text{I}} = \langle exp(\phi) \rangle_{y} = 1$. Looking for a traveling wave, $\phi(y,t) \equiv \phi(y - Ut)$, one can write equation (7) in the form describing motion of a quasi-particle along the coordinate ϕ while y plays the role of effective time

$$d_y^2 \phi = -\kappa^2 + C e^{-\phi} \equiv -\partial_{\phi} W(\phi), \text{ and } (d_y \phi)^2 / 2 + W(\phi) = E$$
(8)

where *C* is the integration constant, $\kappa^2 = \rho_s^{-2}(U_{DW}/U-1)$, $W(\phi) = \kappa^2 \phi + C(e^{-\phi}-1)$, and *E* is the effective energy. Assuming that $\phi(y)$ is periodic, it should be bounded by some 'turning points' $(d\phi/dy = 0, \text{ implying } W(\phi_{max}) = W(\phi_{min}) = E)$ at positive ϕ_{max} and at negative ϕ_{min} (if both ϕ_{max} and ϕ_{min} have the same sign, $\langle \exp(\phi) \rangle_y = 1$ cannot be held). In addition, we should have $W(\phi_{min} < \phi < \phi_{max}) < E$ to ensure that our quasi-particle moves in a potential well, which is only possible when both *C* and κ^2 are positive, which implies $U < U_{DW}$. Our constraint $\langle \exp(\phi) \rangle_y = 1$ can now be expressed in terms of ϕ

$$\int_{\phi_{\min}}^{\phi_{\max}} (\mathbf{e}^{\phi} - 1) \{ \widehat{E} - \widehat{W}(\phi) \}^{-1/2} d\phi = 0, \tag{9}$$

where $\widehat{W}(\phi) = \phi + \widehat{C}(e^{-\phi} - 1)$ and $\widehat{E} = E/\kappa^2$, $\widehat{C} = C/\kappa^2$. The analysis of equation (9) shows [29] that the solution only exists for $\widehat{C} < 1$. The solutions for $\widehat{C} < 1$ correspond to linear drift waves, while for $\widehat{C} \ll 1$ they exhibit strongly nonlinear features and large deviation of the plasma density within the blob, n_{blob} , from the average one. Analytic estimates give



Figure 3. Relative plasma density in the blob (a), $\hat{\delta}_b$ and \hat{L}_2 (b) versus \hat{C} found from numerical solution.

 $n_{\text{blob}}/\langle n_{\text{e}} \rangle \approx 2\sqrt{-(\alpha/\pi)\ln(\widehat{C})} \gg 1$, where $\alpha \sim 1$. The ratio $n_{\text{blob}}/\langle n_{\text{e}} \rangle$, the spatial size of the blob, $\hat{\delta}_{\text{b}}$, defined as the distance between the points with $\phi = 0$, and the distance between the two consecutive blobs, \widehat{L}_2 , (both normalized to κ), following from numerical solution of equation (7), are shown in figure 3 for different \widehat{C} .

Thus, by keeping nonlinear Boltzmann factor in the electron density dependence on the electrostatic potential we demonstrated the existence of large plasma density structures, somewhat similar to those observed in experiments, within the simplest framework of drift wave dynamics.

6. Dynamics of blobs in a sheared magnetic field

As we mentioned in section 5, dynamics of the intermittent filamentary plasma structures, such as blobs and ELM filaments, is an important physical component determining enhanced turbulent transport of particles and energy observed at the plasma edge in fusion devices.

The problems related to the analysis of isolated blob dynamics in the edge plasma were addressed in numerous publications (see e.g. review [24]). However, few papers only were devoted to investigation of blob dynamics in a sheared magnetic field relevant for the tokamaks [31–35]. Here we review the results of the analysis of macroscopic motion of the sheath-connected blobs in magnetic fields having arbitrary geometry of the field lines within the electrostatic approximation. More details of our studies can be found in [36].

For the present study we employ the electrostatic model of blob dynamics. We assume that a sheath-connected blob with the length *L* and the transverse size δ is immersed into the homogeneous plasma background with isothermal electrons having the temperature T_e and with singly charged cold ions having $T_i \ll T_e$. We consider the plasma with low beta, $\beta = 8nT_e/B^2 < m_e/M$ (m_e is the electron mass), so that we can neglect the electromagnetic effects related to propagation of the Alfven waves. In addition, we assume that the magnetic field skin depth $\delta_m = \sqrt{2D_m\tau}$ (where $D_m = c^2/(4\pi\sigma)$ is the magnetic diffusion coefficient, σ is the electrical conductivity of plasma and τ is the characteristic time of blob advection across the magnetic field lines) is large compared to the transverse blob size, $\delta_m > \delta$. In this case, we can neglect the frozenness condition for the magnetic field and assume that the blobs can propagate freely across the magnetic field lines. Finally, we assume that the longitudinal density scales for a blob are large enough to neglect the parallel ion dynamics and plasma redistribution along the magnetic field lines. Detailed analysis of the assumptions employed for this study and their range of validity is given in [36].

The transport equations describing macroscopic motion of the blob include the density and vorticity equations, which are as follows:

$$\mathbf{d}_{\mathbf{t}} n = -(\vec{b} \times \vec{g}) \cdot \nabla n / \Omega_{\mathbf{i}} + \nabla_{||} j_{||} / e, \tag{10}$$

$$\rho_{\rm s}^2 {\rm d}_{\rm t} \varpi = -(\vec{b} \times \vec{g}) \cdot \nabla n / \Omega_i + \nabla_{||} j_{||} / e, \qquad (11)$$

where $\varpi = n \nabla_{\perp}^2 \phi + \nabla_{\perp} n \cdot \nabla_{\perp} \phi$ is the plasma vorticity, \vec{g} is the effective gravity acceleration assumed for the present study to be related only to the curvature of the magnetic field lines, $\vec{g} = -2C_s^2(\vec{b} \cdot \nabla)\vec{b}$, $\vec{b} = \vec{B}/B$, C_s and Ω_i are the ion sound speed and cyclotron frequency and $j_{\parallel} = (\sigma T_e/e) \nabla_{\parallel} (\ln(n) - \phi)$ is the parallel plasma current. At the sheath we employ the standard sheath boundary conditions: $j_{\parallel} = \pm enC_s\phi$, whereas for the plasma density and vorticity we set the Neumann boundary conditions. At the side boundaries of the physical domain, the zero-gradient boundary conditions are set for all plasma parameters determining motion of the filament.

Within the approximations adopted, dynamics of the blobs considered in the present study is determined by $E \times B$ advection of the plasma across the magnetic field lines. Therefore, to analyze the macroscopic motion of the filaments, we have to find the distribution of the electrostatic potential in the blob. Analysis of the vorticity equation (11) [36] shows that for the blobs having the collisional parameter $\Lambda = (L/\lambda_e)(m_e/m_i)^{1/2} \ll 1$, the normalized electrostatic potential ϕ established in the plasma can be represented as a composition of three different components, each being attributed to a different physical mechanism, i.e.

$$\phi = \phi_{\rm g} + \phi_{\rm B,eff} + \phi_{\rm pol}.$$
 (12)

Here $\phi_{\rm g}$ is the regular dipole potential, arising from charge separation in the effective gravity field. $\phi_{\rm B,eff} = \ln(n/n_{\rm eff})$ is the effective Boltzmann potential, $n_{\rm eff} = n_+^{\alpha} n_-^{1-\alpha}$, $\alpha = n_+/(n_+ + n_-)$ and $n_{\pm} = n|_{\pm L/2}$; its formation is driven by inhomogeneity of the plasma distribution along the magnetic field lines, parallel electron dynamics and also asymmetries in the blob plasma distribution at the sheaths with respect to each other. Finally, $\phi_{\rm pol}$ is the potential component related to ion inertia and polarization currents flowing in the plasma. For the case where both the curvature radius, $R_{\rm c}$, of the magnetic field lines the blob size δ are large enough $(R_{\rm c} > L(\rho_{\rm s}/\delta),$ $\delta \gg \delta_{\rm pol} = \rho_{\rm s} (L/\rho_{\rm s})^{1/4})$, the contribution of both $\phi_{\rm g}$ and $\phi_{\rm pol}$ can be neglected compared to $\phi_{\rm B,eff}$ [36]. For this case after some algebra we arrive at the following form of equation (10)

$$\partial_t n = -C_{\rm s} \rho_{\rm s}(\vec{b} \times \nabla \Phi_0) \cdot \nabla n, \tag{13}$$

where $\Phi_0 = -\ln n_{\text{eff}}$. It is possible to show that for the case where the magnetic field metrics at the sheaths are equal,



Figure 4. Density slices of the small blob taken at z = -L/2 (filled contours) and z = L/2 (solid lines) cross-sections of the filament, as observed in the field-following coordinate system.

 Φ_0 does not depend on time and only plasma densities at the sheaths in the initial moment of time determine the blob motion [36].

Next, we consider dynamics of the plasma filaments in a magnetic field with a small curvature assuming equal plasma densities at the sheaths. For small- and mediumsize plasma filaments having the transverse dimension $\delta \leq \delta_*$, where $\delta_* = \rho_s [gL^2/(4C_s^2 \rho_s)]^{1/5}$ is the critical size of the blobs [24], the dominant source of potential dissipation is related to ion inertia so that equation (11) can be reduced to $\rho_s^2 d_t \varpi = -\vec{b} \times \vec{g} \cdot \nabla n / \Omega_i$. Therefore, the distribution of the plasma potential, established in a small blob, is formed in each locally perpendicular to the magnetic field cross-section of the filament independently of the neighboring crosssections. In sheared magnetic configurations this will lead to the blob cross-sections moving with macroscopic velocities having both radial and poloidal components varying along the magnetic field lines. If, in addition, the values of the magnetic shear are disparate at the sheaths, the blob ends will start sliding in different directions, causing formation of the effective Boltzmann potential, which will alter the overall dynamics of the filament. In contrast, for a large blob initially homogeneous along the magnetic field lines, the gravity-driven potential dipole established in each filament cross-section is equal along the whole magnetic field line implying that the potential distribution in the blob is integral rather than local. Therefore, a large blob can be expected to move as the whole filament.

To illustrate the difference between the dynamics of small and large blobs we used the BOUT++ code [37]. The details of simulation parameters can be found in [36]. The simulations results are shown in figures 4 and 5. As one can see, the cross-sections at z = L/2 and z = -L/2 of the small blob get misaligned with respect to each other as the filament propagates radially outwards indicating that different parts of the blob move virtually independently of each other, whereas the large blob shows coherency in the course of its motion.

7. Dynamics of He bubbles in W lattice

It is well known that due to its low solubility, He tends to precipitate into bubbles when embedded in metals. Layers of He nano-bubbles of the thickness $\Delta_b \sim 30-50$ nm were observed experimentally in the near-surface region of He-irradiated tungsten, for He energies below the sputtering threshold and fluxes relevant to the ITER conditions, $\Gamma_{\text{He}} \sim 10^{19} \text{ cm}^{-2} \text{ s}^{-1}$



Figure 5. Density slices of the large blob taken at z = -L/2 (filled contours) and z = L/2 (solid lines) cross-sections of the filament, as observed in the field-following coordinate system.

[38–41]. When He fluence, Φ_{He} , exceeded some critical value $(\Phi_{\text{He}})_{\text{crit}} \sim 2-3 \cdot 10^{20} \text{ cm}^{-2}$, which did not depend on Γ_{He} within few orders of its magnitude [38, 41], the strong modification of surface morphology was initiated. The size of the bubbles in the layer was rather homogeneous over the entire layer and was about a few nm for the temperatures below 1000 K [42].

Theoretical studies of the processes associated with the irradiation of W with He (see [43, 44] and the references therein) are usually based on both molecular dynamics (MD) simulations and cluster dynamics (CD) models assuming that He bubble nucleation proceeds by self-trapping or by trapping in existing immobile traps (e.g. impurities). However, the results following from such models are not quite consistent with the experimental observations. First, for the bubble nucleation via He self-trapping, theoretical estimates of the nano-bubble layer thickness and of the He fluence, at which this layer saturates and one should expect to observe a strong modification of the surface morphology, depend on Γ_{He} , contradicting the experimental data [38, 41]. Besides this, for both the mechanisms of bubble nucleation in a thick sample considered in [44], the size of the nano-bubbles is very inhomogeneous across the layer, which is not consistent with the experimentally observed bubble size distributions [40, 42].

Recently, however, it was demonstrated that the formation of the dislocations and vacancies by the growing He bubbles [45] can result in an avalanche effect, when He precipitation on the induced traps leads to growth of more bubbles, furthering creation of new traps [46]. This effect can strongly facilitate the nucleation and growth of nano-bubbles in both thick samples and relatively thin tendrils. The MD simulations [46] of the behavior of helium atoms in tungsten lattice for two cases-the case of the initially perfect monocrystalline W and the case with an over-pressurized helium bubble of diameter ~2nm pre-formed in the W lattice-showed that the concentration of the mobile He clusters within the dislocations surrounding the bubble exceeds that in the initially unperturbed W lattice by an order of magnitude, the concentration of larger clusters being higher by about two orders of magnitude (see figure 2 from [46]). Thus, the large immobile clusters are formed almost exclusively in the lattice imperfections produced by the growing nano-bubbles, serving as effective He trapping locations. The CD model, including the effects of He absorption by nano-bubbles and the generation of new He traps by bubble growth, was developed in [46]. The generation of additional traps was described by fitting

dimensionless parameter γ showing now many new traps are generated per each He atom absorbed by all clusters. It was shown that for $\gamma = 10^{-3}$ this model gives a very good agreement with currently available experimental data for the fluence critical for the surface morphology modification, the thickness of the nano-bubble layer, and the characteristic size of the nano-bubbles.

8. Conclusions

In this brief review we just touched upon different issues related to the complex and multifaceted physics of the edge plasma in fusion devices. Many processes considered in this review exhibit synergistic effects, which, due to the lack of space, were not considered here in depth. More details can be found in our publications and the references therein.

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