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ELMs IN DIII-D HIGH PERFORMANCE DISCHARGES

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ABSTRACT

A new understanding of edge localized modes (ELMs) in tokamak discharges is emerging [P.B. Snyder, et al., Phys. Plasmas, 9, 2037 (2002)], in which the ELM is an essentially ideal magnetohydrodynamic (MHD) instability and the ELM severity is determined by the radial width of the linearly unstable MHD kink modes. A detailed, comparative study of the penetration into the core of the respective linear instabilities in a standard DIII-D ELMing, high confinement mode (H-mode) discharge, with that for two relatively high performance discharges shows that these are also encompassed within the framework of the new model. These instabilities represent the key, limiting factor in extending the high performance of these discharges. In the standard ELMing H-mode, the MHD instabilities are highly localized in the outer few percent flux surfaces and the ELM is benign, causing only a small temporary drop in the energy confinement. In contrast, for both a very high confinement mode (VH-mode) and an H-mode with a broad internal transport barrier (ITB) extending over the entire core and coalesced with the edge transport barrier, the linearly unstable modes penetrate well into the mid radius and the corresponding consequences for global confinement are significantly more severe. The ELM accordingly results in an irreversible loss of the high performance.
I. INTRODUCTION

Since their first identification in early ASDEX High Confinement Mode (H-mode) plasmas [1], ELMs have become ubiquitous in H-mode discharges both tokamaks [1–6], and even stellarators [7] but their nature has been poorly understood until quite recently, with a wide range of causal mechanisms proposed over the past two decades. Early on, in 1985, it was proposed [8] that ELMs are essentially ideal edge kink modes (or peeling modes [9]) driven by the Pfirsch Schluter current from the steep edge pressure gradient in H-mode; this yielded a limiting edge pressure gradient, \( p_{\text{edge}} \), for instability and qualitatively explained most of the features of the ELM known at the time. However, it was later noted [10] that the onset of ELMs quantitatively correlated with the edge pressure gradient, \( p_{\text{edge}} \), locally reaching a value close to the estimate for the infinite n ballooning limit in the simplified, so-called \( s-\alpha \) model, obtained by ignoring more sophisticated effects from cross section shaping and possible local second stability access [11]. The onset also appeared to scale with this prediction [10]. Consequently, it was fairly widely accepted for several years that ELMs are essentially, ideal, or almost ideal, high n ballooning modes.

Nevertheless, with advances in discharge equilibrium reconstructions, it gradually became clear that this model did not always work – the diverted non circular cross sectional geometry of most H-mode discharges yielded quantitatively different predicted limits and, in some cases with apparent second stability access, no local limit at all. A model was later proposed [12], in which kink or peeling modes, driven by edge current density gradients, were responsible, although, without extensions, this model could not account for the presence of a threshold \( p_{\text{edge}} \) for ELMs.

The discovery that the neoclassical bootstrap current associated with local pressure gradients is real and is measurable, [13] yielded a new element to the model proposed in Ref. [8] — that the bootstrap rather than just the Pfirsch Schluter current density could provide destabilization of the peeling modes [14]. This was investigated in a series of numerical studies of large ELMs [14,15] where it was shown that there is generally a threshold in both \( p_{\text{edge}} \) and edge current density, \( j_{\text{edge}} \), or its gradient, \( j'_{\text{edge}} \), for the destabilization of low to intermediate n peeling-like modes. Analytic work [16,17] later confirmed the numerical results by extending the ballooning formalism to include a vacuum and current density gradient. The unstable modes were accordingly labeled “peeling-ballooning” modes, and this nomenclature will be adopted here.
Despite the promising advances, the theory for peeling-ballooning modes as being responsible for the observed ELMs remained qualitative only. One reason for this was the rather extreme sensitivity of the predicted edge stability to details of the plasma cross section and edge profiles, coupled with the rather poor experimental data available to constrain those profiles. This, however, has been partly remedied over the past decade and the cross section and pressure profiles can now be reasonably well diagnosed. With this, it became apparent that in many cases in DIII-D, the pressure gradient profile, \( p' \), could be several times larger than the predicted first regime ballooning limit [18,19]. This discrepancy can be resolved by requiring some form of consistency between the relatively well-diagnosed pressure profile and its associated bootstrap current, and the relatively less well-diagnosed current density profile [18–22]. With this additional constraint on the profiles, much of the remaining sensitivity of the stability can also be removed.

ELMs are generally categorized as Types I, II, or III [23], with Type I referring to the most common type and which appears to be similar in most large tokamaks [6]. Since there is no uniform definition or characterization of Type II and III across devices, only those of Type I are discussed here. A predictive model for Type I ELMs will need to explain several important features; ultimately, a complete model of ELMs will inevitably require an explanation of the Type II and III events as well. The characteristic features include the magnetic fluctuation signatures, the fast, ideal-like growth rates, the generally close proximity to ballooning limits in some cases and apparent violation of the limits or second stability access in others, and the dependence of ELM features on cross sectional shape [18–22], collisionality [19,20,21], and other conditions. Of particular interest here, the model also needs to explain the wide range in size and effect of ELMs under varying conditions.

In conventional H-mode discharges, Type I ELMs can vary from small, frequent events with repetition rates of the order of several kHz and drops in the edge electron temperature, \( T_e \), of a few eV, to less frequent events below 50 Hz but with much larger effects on the edge \( T_e \) of several hundred eV. Although the largest ELMs can have serious deleterious effects in a large device due to the rather large localized energy and particle fluxes on to material surfaces, they are usually benign with respect to their effect on the plasma in the sense that the degraded confinement and lost stored energy during the ELM are rapidly recovered, and the high confinement remains intact until the next ELM.

In high performance discharges, with enhanced high confinement enhancement factor over the low confinement mode (L-mode) scaling, \( H \geq 3 \) (\( H \leq 2 \) for conventional H-mode), the initial high confinement ELM-free period is invariably terminated by a large ELM event. For these Very High Confinement mode (VH-mode) discharges [24], this first ELM, often referred to as the “X event” [25], terminates the high confinement so that the continual increase in stored energy, characteristic of VH-mode, is ended, there is a sudden
drop in the plasma $\beta$, and the confinement enhancement factor drops to $H \approx 2$ as the
discharge reverts to conventional ELMing H-mode [26,27] ($\beta$ here is the usual ratio of
plasma thermal energy to the magnetic energy of the confining vacuum toroidal field).
Subsequent ELMs do not further affect the H-mode level confinement between ELMs but
the VH-mode is very rarely recovered.

A similar sequence occurs for negative central shear (NCS) discharges [19,28] with
an H-mode edge and internal transport barrier (ITB) in DIII-D. During the ELM-free
H-mode phase these NCS discharges also have $H \geq 3$. However, the first ELM event
invariably destroys the high confinement in the core, but leaves the H-mode transport
barrier intact as the discharge again reverts to a conventional ELMing H-mode.

The VH-mode and NCS H-mode discharges are the highest peak performance
discharges in DIII-D [19,27,28], while in JET [25,29], VH-mode was chosen for the high
fusion yield DT experiments [29]. Consequently, it is important to understand the limit
caused by the first large ELM and its effect on the discharge interior. The reduced
performance and its correlation with the first large ELM are well-documented [19,24-28].
While there is sometimes other MHD activity present, this does not appear to correlate
with the irreversible change in confinement. The only exception here, is the $m/n = 1/1$ mode — usually a sawtooth — which often occurs as well in the VH–mode discharges,
in which case the global energy loss is considerably larger [26]. But the large ELM is
always present.

From the developments during the past decade, a new model for ELMs in Tokamak
discharges, in terms of the interplay between edge MHD kink modes of different toroidal
mode numbers $n$, has evolved recently [19,21,22,30–34], and appears to successfully
explain a wide variety of ELM behavior in DIII-D [22,35,36], JT-60U [22,36], Alcator
C-Mod, [37] and ASDEX-U [38]. Within the framework of this model, the ELM relaxation
cycle is explained as resulting from a complex interaction between high $n$ ballooning and
low $n$, peeling-ballooning mode stability. The ELM itself is considered to be triggered by
an ideal MHD instability, driven by the current density and pressure gradients near the
edge of the plasma, and the ELM severity, given by the depth of the energy loss into the
plasma core, is determined by the radial width of the linearly unstable MHD kink modes.
These are working assumptions of the model and the general predictions obtained from
them appear to hold remarkably well. Whereas many additional effects and complications
need to be invoked for a fully quantitative ELM description, these two key assumptions are
the essential basis.

In this paper, a comparative study of the penetration into the core of the respective
linear instabilities in a standard DIII-D ELMing H-mode discharge, with that for two
relatively high performance discharges is described in the context of the new ELM model.
A key feature of this model is that the range of ELM sizes results from the destabilization of different ranges of $n$ and different penetration of the mode into the core; deeper penetration into the core is assumed to result in more catastrophic ELMs. This assumption is investigated by comparing the edge stability of these three ELMing discharges in DIII-D. The following section discusses the experimental discharges considered. The high performance discharges analyzed are a VH-mode with very high confinement in the ELM-free period and a high performance NCS H-mode discharge with the irreducible minimum, neoclassical level transport over almost the entire cross section [19,28]. In both, the high confinement was irreversibly lost at the first large ELM and the discharges reverted to ELMing H-mode with standard H-mode level confinement. They were chosen as suitable examples since good quality equilibrium data was available in each case and since the toroidal mode number was also lower than usual and could be grossly determined from the Mirnov fluctuation data. Otherwise, they were typical of the VH-mode and NCS H-mode discharges in DIII-D. These are contrasted with a conventional H-mode discharge, where the first large ELM was relatively benign. Section II also briefly summarizes the new model for ELMs developed in Refs. [18–22,30–34] and the numerical procedures used in the work, particularly the equilibrium reconstructions, which are a critical element in the analysis.

Following this, Section III considers the stability results for the three cases, showing that the stability results are consistent with the ELM model. This is followed in Section IV by a discussion of the qualifications and reservations that need to be kept in mind when interpreting these results, along with a discussion of the implications. Section V briefly summarizes the results and the conclusions.

A number of the reservations discussed in Sections IV and V are important and worth emphasizing. First, this work does not attempt to either prove or to explain the physics behind the basic assumptions of the ELM model but is intended only to extend the model to include the two most important high performance discharges in DIII-D. All that can be said of the underlying physics basis is that it is plausible and that the model appears to hold up extremely well under a large range of conditions, including those here. Second, the stability calculations for the edge instabilities are sensitive to the equilibrium reconstructions and these are subject to some uncertainty; not all valid reconstructions reproduce the observed instabilities, especially for the conventional H-mode. The important point is that it is possible to find equilibria consistent with both the discharge data and the ELM model.
II. ELMS IN HIGH PERFORMANCE DISCHARGES

Figure 1 shows the evolution of a conventional, lower single-null (SN) DIII-D discharge #92001 [39]. The upper traces [Fig. 1(a)] are the power injected via neutral beams, the electron density, $n_e$, and the photo diode $D_\alpha$ signal from early in the discharge, through the H-mode transition at 1610 ms, the ELM-free period to the first Type I ELM at 2140 ms, and the later ELMing phase. For this discharge, the total current is $I = 1.465$ MA, the vacuum field is $B = 1.87$ T, and the $Z_{\text{eff}}$ is 2 in the core, rising to 4 at the top of the pedestal and back to 2 at the edge. The lower trace in Fig. 1(b) shows the global confinement enhancement factor, $H$, over the ITER-89P scaling [40]. At the first ELM, the edge pressure gradient drops by a factor 2.3 and the confinement drops several percent but remains above $H \approx 2$ through the remainder of the discharge. The ELM is observed on the Mirnov system but the toroidal mode number $n$ could not be resolved; the Mirnov system on DIII-D can resolve up to $n \approx 6$. In a few conventional H-mode discharges where the mode number is lower, the Mirnov system typically indicates that the magnetic signal from an ELM is a superposition of many intermediate $n$.

![Fig. 1. Evolution of DIII-D H-mode discharge #92001 (a) injected neutral beam power (MW — dotted line), $n_e$ ($10^{19}$/m$^3$ — dashed line), and $D_\alpha$ signal (arbitrary units — solid line), (b) confinement enhancement factor $H$ over the ITER-89P L-mode confinement scaling.]

The discharge equilibrium was reconstructed using the EFITD code [41] at two times, at 1693 ms early in the ELM-free phase, and at 2075 ms just before the first Type I ELM. The reconstruction utilized the 35 channel Motional Stark Effect (MSE) diagnostic to obtain the field line pitch [42,43], in addition to the Thomson system for the electron pressure and charge exchange recombination (CER) for the ion pressure, with the fast neutral beam (NB) contribution obtained from a slowing down calculation. In addition, a
complete set of magnetic flux loops and poloidal field sensors constrained the plasma boundary. The reconstructed equilibria at the two times are displayed in Fig. 2. The cross section and flux surface geometry are given in Fig. 2(a) for the equilibrium at 2075 ms; at 1693 ms these are not noticeably very different. At 1693 ms, \( \beta = 0.95\% \), and \( \beta_N = 0.75 \), and at 2075, \( \beta = 2.14\% \) and \( \beta_N = 1.67 \). Here, the normalized \( \beta, \beta_N = \beta/(I/aB) \), is the Troyon coefficient [44]; \( I \) is the total current in MA, \( a \) is the minor radius in meters, and \( B \) the vacuum toroidal field in Tesla. The pressure gradient profiles are displayed in Fig. 2(b,c) at 1693 ms and 2075 ms respectively. These are shown as functions of the poloidal flux \( \psi \), normalized between zero at the magnetic axis and 1.0 at the plasma boundary. The key features here are the large gradients near the edge of the plasma. At 2075 ms, this is very highly peaked at the edge. Also, note that the reconstructed pressure profile has finite \( p_{\text{edge}} \), and hence finite current at the separatrix.

The pressure and its gradient are well determined by the available data, especially in the pedestal region where both the pressure profile data and the external flux loop and poloidal field probe magnetic data provide strong constraints. The fitted pressure is shown in Fig. 2(d) with the measured profile data. A comprehensive analysis of the uncertainties in the reconstructed profiles is not possible for several reasons. The EFITD code fits the profiles to the available measured data in a statistical \( \chi^2 \) sense [41] and it is difficult to estimate how the errors are propagated through to error estimates for the actual pressure and current density profiles since systematic errors in the measured data dominate the statistical errors. The ion and Thomson data are measured at different spatial points and mapped to the poloidal flux through the equilibrium fitting process. The points and error bars shown in Fig 2(d) are estimated statistical errors from the Thomson data for the electron pressure only. Nevertheless, a rough estimate for the maximum uncertainty in the pressure gradient in Fig. 2(d) is about 30\%, with the true pressure gradient likely higher than the fit value due to spatial averaging in the Thomson system.

Even with the extensive equilibrium data, there is a significant range in the detailed profiles that can be fit to the data near the edge across the H-mode pressure pedestal since the MSE signals there contain a dominant contribution from the local radial electric field [43]. Figure 3 shows two alternative reconstructions at 2075 ms. The pressure profile is well constrained by the Thomson data but additional information is required to constrain the current density. In Fig. 3(a), the current density profile is fit with a low order spline, and no additional constraints beyond the MSE and external magnetic flux and poloidal field data. A ballooning stability calculation then shows that the H-mode pedestal, peak pressure gradient extends through the calculated unstable region and into the second stable region; if correct, this would indicate that ballooning stability is completely irrelevant. However, by constraining the edge current density profile to be aligned with the neoclassical bootstrap current density, as in Fig. 3(b), an equally good fit to the data is obtained in terms of minimizing the statistical \( \chi^2 \) [41], which is the most appropriate
characterization of the quality of the reconstruction. For the reconstruction in Fig. 3(a), $\chi^2 = 11$ for the MSE data, and for the reconstruction in Fig. 3(b), $\chi^2 = 13$; these are reasonable and are not statistically different. In the second case, the transition to the second regime opens up and consistency with ballooning stability is restored.

Fig. 2. Reconstructed equilibria from DIII-D H-mode discharge #92001 (a) flux surface geometry at 2075 ms with conducting wall used in the stability calculations, (b) pressure gradient and safety factor profiles at 1693 ms, with $p_0' = 0.39$ MPa/(Wb/rad), and (c) at 2075 ms, with $p_0' = 1.36$ MPa/(Wb/rad), and (d) fitted pressure profile (Pa) and measured profile data at 2075 ms. The radial coordinate is normalized poloidal flux.
This apparent gross violation of ballooning instability occurs in many DIII-D H-mode discharges and is resolved by imposing the additional constraint that the current density must be aligned with, and a sizeable fraction of, the theoretically calculated bootstrap current. Moreover, as will be discussed later in Section III, consistency between the intermediate n ideal stability and the onset of the ELMs is then also obtained. Some flexibility still remains in the reconstructions since consistency can be restored, while still yielding good fits to the equilibrium data, by taking any fraction of the bootstrap current between roughly 50% and 100%. There is a good deal of current debate over the applicability of the collisionless bootstrap model used here [45] for the edge. The allowed reduction of up to 50% in the collisionless bootstrap model current density is quite consistent, however, with the reduction expected from a collisional model [46]. Reducing this flexibility experimentally will require a significant improvement in the measurement of the edge current density profile. This is planned for future experiments in DIII-D [47].

The observation that the equilibrium reconstruction requires a current density profile that is roughly aligned with the bootstrap current in many DIII-D mode discharges is a key element that led to the new model for ELMs. This model has been described in some detail.
in several recent papers [18,19,21,22,30–36], but is worth summarizing briefly since it forms the basis of and motivation for the work described here.

In this ELM model, a buildup in $p'_{\text{edge}}$, associated with the improved confinement from H-mode, results in increased current density near the edge as the bootstrap current accordingly builds up. This current density is a two-edge sword; it tends to be further destabilizing for peeling instabilities but enhances access to the second stability regime for ballooning stability. If access is not achieved before $p'_{\text{edge}}$ reaches the first regime limit, a high $n$ edge ballooning mode results; the $n$ is determined by non-ideal Finite Larmor Radius (FLR) or finite orbit effects but is generally in the range $20 \leq n \leq 40$. On the other hand, if access is achieved, $p'_{\text{edge}}$ can continue increasing, second stability access to high $n$ modes continues to improve (i.e., extends over a widening spatial region), and this continues until a lower $n$ kink ballooning mode is destabilized. In extreme cases, this can be $n \approx 2$ to 3, though rarely an $n=1$ mode.

VH-mode discharges in DIII-D and JET [48] have shown considerable promise for obtaining high fusion gains but are limited in duration by the first large, ELM-like X event [25–27,48]. These edge instabilities have many of the signatures of the Type I ELMs in conventional H-mode discharges. However, there are some important distinctions. Unlike in a conventional H-mode, these ELM events irreversibly destroy the high performance. The wide edge pedestal is reduced to that of a conventional H-mode, and a central collapse is often observed with a significantly larger loss in the discharge $\beta$ and stored energy. Figure 4, is a time trace of a typical VH-mode discharge in DIII-D, #75121. Here, $\beta_N$, the energy confinement, $\tau^{\text{energy}}$, normalized to the JET-DIII-D H-mode scaling, $\tau^{\text{JET}/\text{DIII-D}}$, [48], and the $D_\alpha$ signal are shown. Note that the increased baseline in the $D_\alpha$ signal can be attributed to increased particle recycling from the wall after the first ELM dumps energy on the wall. The JET-DIII-D H-Mode scaling is approximately twice the ITER-89P L-mode scaling so that $H \approx 2\tau^{\text{energy}}/\tau^{\text{JET}/\text{DIII-D}}$. In the early VH-mode phase before the X event, signified by the spike in the $D_\alpha$ signal, the stored energy and $\beta_N$ are continually increasing as a result of the input NB heating coupled with high confinement enhancement over the outer half of the plasma. The confinement there is at neoclassical levels and the global enhancement factor $H$ is typically between 3 and 4. However, this rapid increase is suddenly terminated at the X event, the high confinement is lost, and $\beta_N$ and the stored energy collapse. The magnetic precursor to this event indicates a highly non-sinusoidal, toroidally localized mode with a toroidal extent suggesting that $n \approx 3$ is the largest single component [14]. For this discharge, there is also an $n = 1$ mode present which results in an internal collapse and a correspondingly larger loss of total stored energy. While common in VH-Mode discharges in DIII-D, this feature is not universal. The discharge then continues in a conventional ELMing H-mode phase and the subsequent ELM events have all the features of standard Type I ELMs.
Fig. 4. Evolution of DIII-D VH-mode discharge #75121 showing $\beta_N$ (%T-m/MA — solid curve) and $\tau_{\text{energy}}^\text{JET/DIII-D}$, the energy confinement normalized to the JET-DIII-D scaling (dashed curve), and $D_\alpha$ signal (arbitrary units). The L-mode, H-mode, and VH-mode phases are indicated.

The equilibrium reconstruction for this double-null (DN) discharge at 2590 ms, just before the X event, is shown in Fig. 5. For this equilibrium reconstruction, $I = 1.6$ MA, $B = 2.1$ T, $Z_{\text{eff}}$ in the core is 1.1 and at the edge is 1.9, $\beta = 3.1\%$, and $\beta_N = 2.8$. Unlike the conventional H-mode case, there is relatively little ambiguity in the edge profiles since the edge gradient is much broader and reasonably well resolved by both the Thomson and MSE diagnostics. For the reconstruction shown in Fig. 5. Figure 5(c) shows the pressure profile data from a slightly earlier time at 2540 ms. This is shown as a function of the square root of the normalized volume $\theta$. For the pressure fit in this case, $\chi^2 = 18$. For the MSE data, $\chi^2 = 5.8$, and for the external magnetics, $\chi^2 \sim 90$, which is rather high and indicates some larger than normal systematic error. The large pressure pedestal extends inward over a substantial radius (around 25%) and the peak in the edge current density profile is reasonably well resolved by the MSE diagnostic; for this discharge, 8 MSE channels were available. Ballooning stability calculations using the BALLOO code [49] show second stability access in the edge pedestal region of the plasma [48], and the pressure gradient profile is consistent with ballooning stability.

A similar situation is observed in DIII-D H-mode discharges formed from L-mode NCS discharges with an ITB [50]. After the H-mode transition in these discharges, the strong ITB and highly peaked pressure profile are moderated. The ELM-free H-mode phase is then characterized by a broad, edge transport barrier merging with the ITB to produce an equilibrium with essentially neoclassical levels of confinement over a large fraction of the cross section [51] and corresponding enhancement factors in excess of H~3. As in the VH-mode discharges, this high performance is terminated by the first large ELM [19,51] which has almost identical characteristics to the VH-mode X event. Again, the discharge continues as a conventional ELMing H-mode [51]. One significant difference, however, is that, in the high performance phase, these discharges tend to have elevated safety factor profiles, q, with negative or weak shear and $q \geq 1.5$ in the core, in contrast to the VH-mode where $q = 1$ in the core. This difference will be discussed in more
Fig. 5. Reconstructed equilibrium for DIII-D VH-mode discharge #75121 at 2590 ms. (a) Flux surface geometry showing conducting wall used in the stability calculations, (b) safety factor (solid curve), pressure (character “p”), and normalized density profiles (character “n”) as functions of poloidal flux, and (c) fitted pressure profile (Pa) and measured profile data at 2540 ms. The radial coordinate $\vartheta$ is the square root of the normalized volume.
detail in Section IV. The time evolution of one of the highest performance NCS H-mode discharges in DIII-D, #87099, is reproduced in Fig. 6 from Ref. [19]. $\beta_N$ is shown in Fig. 6(a) with the Mirnov signal trace; the first large ELM appears on this trace as a single large spike. For this discharge, $I = 2.0$ MA and $B = 2.05$T. Also, $Z_{\text{eff}} = 2.1$ in the core and 2.8 near the edge. Figure 6(b) shows the edge $T_e$ and the $D_\alpha$ signal. At 1590 ms, the $D_\alpha$ signal drops, indicating an H-mode transition, and the edge $T_e$, measured with approximately a 15% relative error near the top of the pedestal at about $\psi = 0.9$, increases suddenly. This edge $T_e$ then collapses at the ELM and never fully recovers. The total stored energy increases during the L-mode phase due to the strong ITB, but this increase is accelerated after the H-mode transition. During the ELM-free H-mode phase, this discharge had $H \sim 3$ and neoclassical confinement measured over a large fraction of the discharge cross section. After the first ELM, however, only the conventional H-mode barrier remained and $H$ dropped to $H \sim 2$. In this discharge, the magnetic precursor signal on the Mirnov diagnostic was resolved and found to be a highly non-sinusoidal superposition of multiple $n$, roughly around $n \sim 5$–6 [19], indicating that the mode, when observed, was in a strongly nonlinear stage; $n$ was again estimated from the mode extent of about one sixth of a toroidal period. The growth time was estimated at $\gamma^{-1} \sim 150$ ms.

![Fig. 6. Evolution of DIII-D NCS H-mode discharge #87099 (a) $\beta_N$ (%mT/MA) and Mirnov signal (arbitrary units), (b) edge $T_e$ (keV) measured near the top of the pedestal and $D_\alpha$ signal (arbitrary units). The nominal relative error in $T_e$ is 10% to 15%. The NCS H-mode phase is between the two dashed vertical bars.](image)

The reconstructed equilibrium for this DN discharge at 1875 ms is given in Fig. 7. In this case, $\beta$ was increased by a few percent of its best fit value of $\beta = 3.7$%, in order to find unstable finite n modes in the stability calculations with reasonable resolution.

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Fig. 7. Reconstructed equilibrium for DIII-D NCS H-mode discharge #87099 at 1875 ms. (a) Flux surface geometry showing conducting wall used in the stability calculations, (b) safety factor (solid curve), pressure (dashed curve), and density (short dash curve) profiles, and (c) fitted pressure profile (Pa) and measured profile data at 2075 ms. The radial coordinate $\theta$ is the square root of the normalized volume.

However, this increase was well within the uncertainty of the equilibrium measurements and $\chi^2$ was only slightly increased. The unscaled fitted pressure is shown in Fig. 7(c) with the measured profile data. The scaling in $\beta$ was applied by simply increasing the pressure
data points uniformly by 10%. At this time $\beta_N = 2.4$. As with the VH-mode discharge, the pedestal region is very broad and well resolved by the MSE system so there is little ambiguity in both the equilibrium profiles –15 radial channels were available in this case and the statistical $\chi^2 = 14.1$ for the fit to the MSE data, $\chi^2 = 5$ for the fit to the pressure data, and $\chi^2 = 60.5$ for the external magnetic data, which is reasonable in this case. Also, ballooning stability calculations indicate that the outermost 7% of the flux is in the transition regime for second stability and, inside that, is just below the first regime limit [19]. The core is also typically in the second stability region for these NCS H-mode discharges.

In both types of high performance discharge, there is little difference in the magnetic and $D_\alpha$ signatures between the first ELM event and the subsequent ELMs, except for an increased magnitude. The major difference is simply the sudden loss of stored energy and collapse of the VH-mode or NCS H-mode internal transport barrier. An understanding of the differences between the X events and the Type I ELMs is essential for sustaining the VH-mode or NCS H-mode discharge performance and realizing the full potential of their enhanced confinement.

Some detailed data on the profile changes due to the ELM is available from analysis of a number of conventional H-mode discharges. Most noticeably, the $p'_{\text{edge}}$ typically drops significantly; the pressure gradient is reduced by a factor of 2.3 in discharge #92001 but up to 3 for larger ELMs and the reduction varies considerably along with the total stored energy loss [52]. For the conventional H-mode, the profiles are modified into about $\psi = 0.8$. For the high performance cases, the affected region is much greater, in to $\psi = 0.6$ or more [26].

For the equilibria reconstructed from the conventional H-mode discharge #92001, the VH-mode discharge #75121, and the NCS H-mode discharge #87099, low and intermediate $n$ stability up to $n = 5$ was evaluated using the GATO ideal MHD stability code [53]; higher $n$ requires prohibitively large resolution to resolve the modes. For edge stability with intermediate $n$ even in this range, considerable care was required to resolve the unstable modes. This will be discussed in some detail in Section IV. Some higher $n$ modes were also studied with the ELITE code [54], which is valid for $n \geq 10$. Analysis in the range $5 < n < 10$ is difficult but the results in the lower and higher ranges are consistent and there is no reason to expect different results in this intermediate range. It is also important to emphasize that the bootstrap alignment constraint is imposed in the conventional H-mode case, but is already clear and present in the resolved data for the two high performance discharges, so is not imposed as an additional constraint. Without the bootstrap current near the edge, it is generally not possible to find low $n$ instability that correlates with the observed ELM, even with MSE data, in the conventional H-mode. Previous attempts to find such a correlation in actual discharges have generally failed [55].
Some limited sensitivity studies were performed with varying equilibrium profiles to determine the sensitivity of the results. In general, it is not possible to claim that the discharges are unambiguously unstable to the edge modes obtained here; for example, for the conventional H-mode discharge #92001, the profiles shown in Figs. 3(a) and 3(b) are equally good fits to the measured data. The implications of this will be elaborated further in Section IV. However, for the VH-mode discharge #75121, a more extensive set of sensitivity studies was performed to obtain a better understanding of the physics. This will be discussed in the following section.
III. LOW AND INTERMEDIATE N EDGE STABILITY

A. VH-mode

For the VH-mode discharge #75121, the low to intermediate n linear stability was computed for the equilibrium reconstructed from the discharge data at 2590 ms. Stability to n=1 edge modes was found both with and without a wall, although an internal n =1 kink mode was found to be unstable, consistent with the equilibrium having an axis safety factor value of $q_0 \sim 1$. For n=2, 3, and 4, however, an instability was found even with the DIII-D wall. This is shown in Fig. 8 for the n=3 mode. The mode is clearly localized near the edge of the plasma and the mode structure is qualitatively consistent with the general features expected for an X event in several respects. The edge localization is easily seen in the plots of the fluid displacement vector $\mathbf{\xi}$, projected onto the poloidal plane in Fig. 8(a). Although peaked towards the edge, the mode clearly penetrates well into the core of the plasma. This is most clearly seen in the Fourier decomposition of the normal displacement, $X = \mathbf{\xi} \cdot \mathbf{\hat{\psi}}$ in Fig. 8(b). Note, for example, that at $q=2$, corresponding to $\rho \approx 0.75$, the mode amplitude is still considerable. Here, $\rho$ is a normalized radius obtained from mapping the poloidal flux on the outboard side to the spatial midplane distance. This is a convenient mapping from $\psi$ since it shows the penetration of the mode in real space more accurately. The $\psi$ coordinate is marked at the top at selected rational q surfaces. The Fourier analysis is with respect to the PEST straight field line poloidal angle [56]; the actual stability calculations, however, utilized an equal arc length poloidal angle [53]. On the outboard side, the individual harmonics add constructively, yielding the strong ballooning character seen in Fig. 8(a).

Extensive sensitivity studies were performed using this discharge as a base while varying the magnitude of the edge pressure gradient, the edge current density parameterized by the ratio of current at 95% flux, $J_{95}$, to the average over all flux surfaces ($\mathcal{I}$), and the internal inductivity $\ell_i$. The results from these studies are summarized in Fig. 9. The termination X-event was identified previously with edge peeling-ballooning modes driven by the combination of pressure gradient and current density in Refs. [14,15], and results, in agreement with those shown in Fig. 9, were reported in Ref. [15]. Although the growth rate changes with the magnitude of $J_{95}$, the mode structure is not much affected.

The dependence of the low n mode growth rates on $p'_\text{edge}$ at fixed edge $J_{95}/(\mathcal{I})$ and fixed $\ell_i$ is seen in Fig. 9(a). The equilibria were kept at constant $\beta$ by adjusting the central pressure gradient as $p'_\text{edge}$ is varied. At low $p'_\text{edge}$ therefore, internal modes driven by the
Fig. 8. Computed peeling-ballooning instability for n=3 in the equilibrium for DIII-D VH-mode discharge #75121 at 2590 ms. (a) Projection onto the poloidal plane of the displacement vectors $\mathbf{x}$, (b) Fourier decomposition of $X = \xi \cdot \nabla \psi$ as a function of the poloidal flux mapped to the radial outboard midplane normalized radius. Poloidal flux values are indicated for selected rational $q$ values.
Fig. 9. Sensitivity study for peeling-ballooning instabilities in VH-mode equilibria based on DIII-D discharge #75121 showing computed growth rates versus (a) $p'_{\text{edge}}$, (b) $J_{95}/\langle J \rangle$, and (c) $\ell_i$ with the other parameters fixed as shown in each case.
high internal pressure gradient were destabilized. At high $p'_{\text{edge}}$ however, in the observed range for VH-mode discharges, $n = 2$ and 3 modes similar to those in Fig. 8 were destabilized. In this range, the mode structure is also not much affected by the magnitude of $p'_{\text{edge}}$. This instability has a threshold in $p'_{\text{edge}}$, for this particular set offixed values of $J_{95}/\langle I \rangle = 0.6$ and $\ell_i = 1.2$, of $p'_{\text{edge}} \approx 1.1$ MPa, which agrees well with the observed threshold values in discharge #75121 and similar VH-mode discharges. Similarly, at fixed $p'_{\text{edge}}$ and $\ell_i$, the peeling-ballooning modes are destabilized above a threshold in $J_{95}/\langle I \rangle$ as shown in Fig. 9(b). Again, the $n=1$ mode is stable and $n = 2$ and 3 become unstable as $J_{95}/\langle I \rangle$ increases, with $n=3$ appearing just before $n=2$.

The dependence on $\ell_i$ is more complicated. Above the thresholds in $p'_{\text{edge}}$ and $J_{95}/\langle I \rangle$, varying $\ell_i$ over a considerable range, yields unstable modes [Fig. 9(c)] with changing character. At high-$\ell_i$ these are largely internal modes and irrelevant to ELMs. At moderate-$\ell_i$, they are peeling-ballooning modes but the peeling component increasingly dominates as $\ell_i$ is decreased. At low-$\ell_i$, the instabilities are localized peeling modes and in between there is a range of $\ell_i$ with low or vanishing growth rates. This study was restricted to $n \leq 3$ but the results from the study of the reconstructed discharge equilibrium indicate that higher $n$ modes up to $n= 4$ are increasingly unstable. Results at even higher $n$, in the range $10 \leq n \leq 40$, obtained using the ELITE code for both conventional H-mode and NCS H-mode equilibria give no indication that the trend does not continue.

These studies show clearly the result, first reported in Refs. [14,15], that the peeling-ballooning modes that are found to correlate with the onset of the large Type I ELMs, are driven by the combination of edge pressure gradient and current density. In VH-mode discharges, this current density is well resolved and clearly present in the equilibria reconstructed from the experiment. It is also consistent with the expected bootstrap current and is responsible for the access to second stability for high $n$ ballooning that is typically found in these discharges [26,29,48].

B. NCS H-mode

For the equilibrium reconstructed from the NCS H-mode discharge #87099 at 1875 ms, stability was found for $n \leq 2$, and $n = 4,5$ modes were calculated to be unstable; the $n = 3$ mode was marginally unstable as determined by the mesh convergence study discussed in Section IV. These were unstable both with and without the nearby DIII-D vacuum vessel wall. The unstable $n=5$ peeling-ballooning mode is shown in Fig. 10. Overall, this mode is similar to that in Fig. 8; the mode in Fig. 10(a) is strongly ballooning on the outboard side (individual harmonics in Fig. 10(b) have the same phase), and is strongly peaked toward the edge but extends well into the half radius. The $n=4$ mode is similar with roughly the same radial envelope. From ELITE
Fig. 10. Computed peeling-ballooning instability for $n=5$ in the equilibrium for DIII-D NCS H-mode equilibrium #87099 at 1875 ms. (a) Projection onto the poloidal plane of the displacement vectors $x$, (b) Fourier decomposition of $X = \xi \cdot \nabla \psi$ as a function of the poloidal flux mapped to the radial outboard midplane normalized radius. Poloidal flux values are indicated for selected rational $q$ values.

calculations with $n \geq 10$, the modes become increasingly unstable with increasing $n$, asymptotically increasing linearly with $n$. The result that a large range in $n$ above $n = 3$ is unstable is consistent with the observation that the precursor to the large ELM consisted of a superposition of modes localized to about one sixth of a toroidal circuit [19].
Limited sensitivity studies were performed by increasing the pressure profile uniformly in increments of 5% of its value, while continuing to fit the remaining equilibrium data as well as possible. The equilibrium of Fig 7 used in the stability calculations has $b$ increased by 5%. Up to about a 10% or 15% increase, this still provided reasonable, although steadily degraded fits — all are well-converged equilibria in the sense of satisfying force balance to high accuracy, however. The stability calculations for each individual $n$ then found increasingly more robust instabilities — more easily resolved and with increasing growth rates — consistent with the results in Fig. 9(a) for the VH-mode discharge, and indicated that the original reconstructed equilibrium of Fig. 7 is only slightly above marginal stability.

This equilibrium has a slightly inverted $q$ profile with $q \geq 2$ over a broad internal region. Hence, its internal magnetic structure is quite different from the VH-mode in Fig. 5. Yet the edge pressure gradient pedestal is similar, and this is what drives the instability in both cases — directly and through the bootstrap current. The differing magnetic structures yield edge instabilities in Figs. 8 and 10 with a different harmonic mix but with roughly the same envelope.

C. Conventional H-mode

For the H-mode discharge #92001 early in the ELM-free period at 1693 ms, the low $n$ stability analysis showed the reconstructed equilibrium to be stable for $n=2$ to $n=5$, even with no wall [35]. For $n=1$, the equilibrium was unstable to an $m=1$ quasi-interchange mode [57]. The best fit to the equilibrium yielded a flat $q$ profile in the core, with a value at the magnetic axis of $q_o = 1.02 \pm 0.1$. Raising $q_o$ to 1.05 by a slight increase in the vacuum toroidal field eliminated the instability. This is well within the experimental errors. If $q_o < 1$ — also within the experimental errors — a conventional internal kink is found. Since this discharge was sawtoothing with $q_o \sim 1$, the stability analysis is then consistent with the observations. Significantly, there is no instability that could be associated with an ELM and no ELM was observed in the experiment.

At the later time, 2075 ms, just before the observed large ELM, the stability calculations showed $n=1$ and $n=2$ modes to be linearly stable, again both with and without a wall. For this equilibrium, $q_o = 1.13 \pm 0.1$. However, instability for $n=3,4,$ and 5 was found for the equilibrium with the DIII-D vacuum vessel; the instability found without a wall is almost identical. In each case, the instabilities are peeling-ballooning modes strongly localized near the plasma edge. Figure 11 shows the $n=3$ mode, with the displacement projection on the poloidal plane in Fig. 11(a) and the Fourier decomposition of $X = \xi \cdot \nabla \psi$ in Fig. 11(b). The $n=5$ instability is displayed in Fig. 12. Note, in particular, the similar radial envelope over the poloidal decomposition between the $n=3$ and $n=5$
modes, with only a slight decrease in the mode penetration for higher n; the n=4 is intermediate between these.

Fig. 11. Computed peeling-ballooning instability for n=3 in the equilibrium for DIII-D conventional H-mode discharge #92001 at 2075 ms. (a) Projection onto the poloidal plane of the displacement vectors $x$, (b) Fourier decomposition of $X = \xi \cdot \nabla \psi$ as a function of the poloidal flux mapped to the radial outboard midplane normalized radius. Poloidal flux values are indicated for selected rational $q$ values.
Fig. 12. Computed peeling-ballooning instability for n=5 in the equilibrium for DIII-D conventional H-mode discharge #92001 at 2075 ms. (a) Projection onto the poloidal plane of the displacement vectors $\mathbf{x}$, (b) Fourier decomposition of $X = \mathbf{\xi} \bullet \nabla \psi$ as a function of the poloidal flux mapped to the radial outboard midplane normalized radius. Poloidal flux values are indicated for selected rational q values.

Higher n modes were analyzed for this discharge using the ELITE code [54] and show little difference in overall structure beyond the continual slight decrease in mode penetration with increasing n. This is shown in Fig. 13. Here, the envelope of the modes are displayed as a function of $\psi$ for the n = 3 mode from GATO and the n = 15, 25, and 40 modes from ELITE. Here, the overall mode amplitudes, which are arbitrary in both codes, were renormalized to match the peak amplitudes. The change in penetration is only moderate for this large range of n.
Fig. 13. Envelopes for the Fourier decomposition of $X = \vec{\xi} \cdot \nabla \psi$ for the computed peeling-ballooning instability for $n = 3, 15, 25,$ and $40$ in the equilibrium for DIII-D conventional H-mode discharge $\#92001$ at 2075 ms showing the slightly decreased penetration in poloidal flux $\psi$ with increasing $n$.

Figure 11 should be contrasted directly with Fig. 8 for the $n=3$ instability in the VH-mode discharge and Fig. 12 with the $n=5$ instability in the NCS H-mode discharge of Fig. 10. In both cases, the conventional H-mode peeling-ballooning instability is much more strongly localized in the edge region. The envelope amplitude is more than 10% of the edge value only in the outermost radial 10% in Figs. 11 and 12 but for nearly half of the radius in Figs. 8 and 10. The edge peaking is even more pronounced if the individual harmonics are summed across the outboard mid plane.

Note that for each of the reconstructed discharge equilibria, the modes shown are fully compressible and were computed using the measured density profile in order to obtain the correct mode structure; both incompressibility and a constant density profile can modify the mode structure (and growth rate), even though the marginal stability point does not change. However, the modifications to the detailed mode structures are small and the overall features are hardly changed. Also, the arbitrary zero in toroidal angle is chosen in the stability calculations to maximize the real part in the Fourier decomposition plots. For the almost up-down symmetric VH-mode and NCS H-mode equilibria, this makes little difference since the real and imaginary parts of the decomposition have almost identical form. For the up-down asymmetric equilibrium in Figs. 11 and 12, the envelope and harmonic mix changes somewhat with toroidal angle (the real and imaginary parts are different and add in varying linear combinations with changing toroidal phase) but the peaking near the edge remains much the same at all angles.
IV. DISCUSSION

Attempts to search for ideal intermediate $n$ kink instabilities in reconstructed conventional or high performance H-mode discharge equilibria prior to 1994 generally failed to find a mode that correlated with observed ELMs. Several critical improvements in the equilibrium reconstructions have, however, since made it possible. Foremost among these are routine, temporally and spatially resolved measurements of the pressure profile and field line pitch, coupled with extensive external magnetic measurements. This enabled the X event terminations in VH-mode to be identified in 1994 [14,15] as the peeling-ballooning modes discussed here. However, although the X event resembled a large Type I ELM, the identification of the peeling-ballooning instabilities with Type I ELMs for specific, conventional H-mode discharges with much more localized pressure pedestals, required several other developments in the modeling.

One element was the improved fitting of the rapid spatial variations in the edge pressure using either hyperbolic tangent representations [18], or spline fits with a sufficient number of knots near the edge. The real key, however, was the realization that the large, localized edge pressure gradients that are then obtained, are completely inconsistent with high $n$ ballooning stability unless the local current density in the edge region was assumed to be aligned with the bootstrap current. In that case, consistency is restored as a transition to second stability opens up. This alignment was suggested by the reconstructions for VH-mode discharges, which showed a broad current density peak associated with the broad pressure pedestal. It is an additional physically plausible assumption that is quite consistent with the internal MSE and external magnetic data, but is not required by the data in the conventional H-mode case.

Then, with sufficient attention to the details in the edge geometry, using the actual reconstructed boundary, profiles, and conducting wall, and with sufficient resolution, low to intermediate $n$ ideal peeling-ballooning modes are found that correlate with the observed ELMs.

Several points should be noted in this respect. First, in each case, the equilibria are up-down asymmetric and diverted. These geometric features need to be respected in the stability calculations. Symmetrizing the cross section to reduce computational time often leads to stabilization of the instability. Similarly, it is not possible to obtain a consistent result if the separatrix boundary is replaced by a plasma boundary moved inward to avoid the X point, by cutting off the edge of the plasma. None of the calculations reported here
were cut off; each used the actual boundary provided from the EFITD reconstruction. Sensitivity studies showed that the results were insensitive to a cut-off of the order of 0.1% of the poloidal flux or less. Otherwise, with a cut-off to any point inside about 99.5% of the flux, the stability is sensitive to the presence or absence of relatively low order rational surfaces just inside the boundary; the stability oscillates with the cut-off and the corresponding instabilities are dominated by pure peeling modes associated with limited plasma boundaries [9].

Second, the stability calculations assume a realistic model for the conducting vacuum vessel in order to properly account for its stabilizing influence on the edge instabilities of interest. A model for the actual vessel wall shape, represented by splines through a set of 60 points, is used, but with rounded corners to avoid numerical problems with the Greens function formulation for the vacuum fields. Two versions are used. For the earlier discharges, including #75121 and #87099, the wall was nearly up-down symmetric and the model used is shown in Figs. 5(a) and 7(a). In the later discharges, such as #92001 however, the wall was modified by the addition of an upper divertor baffle and this is represented in Fig. 2(a). Note that neither wall is particularly close to conformal — especially that in Fig. 2(a). While the presence or absence of a wall and the wall shape can affect whether a given equilibrium is stable or unstable, it should be emphasized that sensitivity studies show that the change in the wall is not responsible for the changes in the calculated mode structures.

Third, high resolution in both radial and poloidal angle grids is generally necessary. This was achieved through a combination of strong grid packing near the edge and the use of mesh convergence studies at each n to ensure adequate resolution and extrapolate the growth rates to zero mesh size. For the more localized instabilities in the conventional H-mode discharge, this is necessary in order to confirm the stability result obtained at moderate mesh values; typically $N_\psi \times N_\chi \sim 100 \times 200$ flux surfaces and poloidal angles is considered moderate, and the results remain consistent up to the finest meshes at $N_\psi \times N_\chi = 300 \times 600$. In Fig. 14(a), the computed growth rates for $n=1$ through 5 in discharge #92001 are plotted against the inverse of $N_\psi \times N_\chi$; a straight line in this figure represents the expected asymptotic quadratic convergence. The only exception to this asymptotic quadratic convergence is that the $n=2$ mode appeared to actually be marginally unstable at the highest resolution, suggesting that the finest meshes are barely sufficient to be in the asymptotic range to properly resolve that case. But this exception is highly sensitive to small changes in the equilibrium and can probably be safely ignored since the growth rates are also extremely small. In all the other cases in Fig. 14(a), the results are well converged.
Fig. 14. Convergence of growth rate with respect to mesh $N_\psi \times N_\chi$ for (a) H-mode equilibrium #92001, showing $n=1$ through 5, (b) VH-mode equilibrium #75121 at 2095 ms showing $n=1$ through 4, and (c) NCS H-mode equilibrium #87099 at 1875 ms showing $n=1$ through 5. The growth rates shown in (b) and (c) were fully compressible calculations.

Figure 14(b) shows the convergence plot for the $n=1$ through 4 modes in the VH-mode case. Here, with coarse meshes, the GATO code finds the usual numerically destabilized Alfvén continuum modes [58] to be the most unstable. The physically unstable mode emerges above the continuum for $n = 3, 4,$ and 5 once the mesh is sufficiently fine, and then converges quadratically, though fairly steeply, to a finite growth rate; the quadratic convergence indicates that the mesh is fine enough to resolve the mode. For $n=1$ and $n=2$ no physically unstable mode is found even with the finest meshes.
For the NCS H-mode discharge, #87099, the instability required a surprisingly fine mesh to resolve at all, despite the fact that the physical instabilities are much less localized than in discharge #92001 and the fact that the required resolution for the VH-mode case was generally much less. The reason for this is not entirely clear but appears to be due to the relatively high q values, which thus requires correspondingly better poloidal resolution since locally, \( m \approx nq \). Experience suggests that insufficient poloidal resolution tends to be numerically stabilizing in GATO even though the Finite Hybrid Element Method [58] employed is normally numerically destabilizing. The convergence plots for the \( n = 1 \) through 5 modes in this case are given in Fig. 14(c). Again, with coarse and moderate meshes, the GATO code finds only the numerically destabilized Alfvén continuum modes but the physically unstable \( n = 4 \) and 5 modes emerge above the continuum once the radial mesh is sufficiently fine to resolve the continuum modes and the poloidal mesh is sufficient to resolve the physical mode. The computed growth rates then converge steeply but quadratically, to a finite growth rate. For \( n = 3 \), the physical mode only emerges with the finest meshes. Indications are that the mode is being resolved but the extrapolated growth rate is small indicating it is quite close to marginal. The \( n = 1 \) and 2 modes again are stable.

Several other cautionary remarks need to be made regarding these calculations. First, and most important, the stability results are quite sensitive to variations in the reconstructed edge current density and since there is often significant ambiguity in the equilibrium reconstruction, especially for conventional H-mode equilibria, the peeling ballooning modes are not always found to be unstable right before an ELM. For example, if the reconstruction from Fig 3(a) is used, no instabilities are found except very high \( n \) ballooning modes that have no correlation with the ELM. The requirement of “reasonable alignment” of \( j_{\text{edge}} \) with the bootstrap current does not fully constrain \( j_{\text{edge}} \); for example, there is some choice in which theoretical model to use — whether collisional effects are important or not — and there is the possibility of additional Ohmic current density, at least in some discharges. For the bootstrap current in Fig. 3(b) used to reconstruct the equilibrium that was employed for the calculations in Figs. 11, 12, 13, and 14, the full collisionless model [45] for the bootstrap current was used. Consistency with second stability access for the measured edge pressure gradient profile can be obtained by using as low as 60% of this calculated bootstrap current. Lower edge bootstrap current, however, tends to preferentially stabilize the lower mode number peeling-ballooning modes [30–32], especially \( n=2 \) and 3 so that the lowest \( n \) that is unstable is increased, which is consistent with the fact that for the ELMs in the conventional H-mode discharges, the observed \( n \) is generally centered around higher values than \( n=4 \). Collisional corrections [46] are expected to be of the order of a 10% to 20% reduction, and are in the right direction but a more stringent quantitative comparison than this is not yet possible without a better measurement of the edge current density profile [47]. Since the correct collisionality model
for the bootstrap current in the edge has not been verified experimentally, and with the other significant uncertainties, it is prudent here to utilize the simpler collisionless model and test the sensitivity of the results to the overall magnitude. The important point, however, is that equilibria that are unstable to intermediate n peeling-ballooning modes are consistent with the equilibrium data just before a Type I ELM but not early in the ELM-free phase. For the high performance discharges, this ambiguity is much less an issue, since the current density peak is spread over a larger pedestal region and is reasonably well resolved.

Also, there are a small number of exceptions to the claim that the first ELM or X-event in high performance discharges irreversibly destroys the high confinement. Analysis of these exceptions is ongoing but, despite the approximately similar values of the edge pressure gradient, the stability calculations have failed to find an unstable peeling-ballooning mode up to n=4. It is not yet clear what is the crucial difference in these discharges, but they tend to be more strongly shaped, and this might be a key; it is not yet clear whether this might just be a result of higher q_{95} or it really is due to the change in shape, in which case different types of shape change at the same q_{95} should produce the same result. Most likely it is due to a combination since elongation, triangularity, and squareness generally have different effects on both low and high n stability [59]. Also, q_0 is often higher, although q_{min} is generally in the same range. A more detailed measurement of the edge current profile [47] should also help to resolve this.

With these provisos, the results of this study are consistent with the hypothesis that the termination event in VH-mode and NCS H-mode is essentially a large, partly pressure driven and partly current driven, peeling-ballooning mode that extends significantly further into the core than in conventional H-mode discharges. This, therefore, conforms with the basic assumption of the recent model developed for Type I ELMs — that the seriousness or size of the ELM, as measured by the magnitude and permanency of the effect on the total stored energy, is determined by the extent of the penetration of the linear modes into the core. Thus, the large Type I ELMs in the high performance cases that destroy the high confinement can be naturally incorporated into the model for Type I ELMs developed in Refs. [18–22,30–38].

The linear mode width appears to be strongly correlated with the width of the edge pressure gradient. For the high performance discharges with a high confinement region extending well into the core, the linear mode extends into the mid radius and the resulting ELM is then more dangerous and leads to an irreversible collapse of the high confinement. The width of the large edge pressure gradient is, in turn, determined to a large extent by the width of the region with access to second stability. In the VH-mode and NCS H-mode discharges, this is typically much broader than in the conventional H-mode. Figure 14 shows the time evolution of the calculated penetration of the second stability access,
\[ \psi_2, \text{ into the core for the VH-mode discharge, calculated from the BALOO code [49]. Access increases slightly with time as } \psi_2 \text{ decreases from } \psi_2 = 0.89 \text{ early, to } \rho_2 = 0.86 \text{ just before the first ELM. The profile in relation to the local high n stability boundaries for the NCS H-mode discharge at 1750 ms, slightly before that used for the stability calculations in Fig. 10, is given in Fig. 4(a) of Ref. [19], showing second stability access from } \psi \sim 0.93 \text{ to 1.0. In both cases, a larger fraction of the plasma edge region than in discharge #92001 [Fig. 3(b)] has access. In all three cases, this region correlates well with the pedestal region of large pressure gradient.}

It should be noted that the linear mode width and therefore ELM penetration is not strongly dependent on the mode number, n, although this was originally an hypothesis of the new ELM model [19–21]. A slow decrease in mode width with increasing n was noted in Section III C for the conventional H-mode case and illustrated in Fig. 13. A similar, slight decrease in radial penetration at higher n is also seen for the NCS H-mode case in both the GATO and ELITE calculations. For the VH-mode, this is also consistent for the computed n=3, 4, and 5 instabilities. This weak dependence on n is consistent with the rough constancy of the radial envelope between the VH-mode and NCS H-mode, despite the different poloidal harmonic content, mentioned earlier; the harmonic content similarly changes with n since the dominant poloidal mode number tends to vary radially as \( m \sim n q \). Although higher n does generally produce more localized instabilities, it is not the prime determinant by itself. Lower mode numbers are fairly well correlated with broader pedestals and larger ELMs but it does not seem to be a direct cause.

Access to second stability, however, depends strongly on cross section shape [19,20,59]. Thus, there is some hope that the virulence of the first Type I ELM in these discharges, might be moderated by varying the discharge cross section. This possibility remains to be fully explored but experiments have already shown that changes in higher order cross section shaping factors can greatly modify the ELM size and frequency [21].

Comparison of the VH-mode and NCS H-mode results is also instructive. As mentioned earlier, these have significantly different q profiles, but the linearly unstable modes and consequent ELMs are essentially little different. This is surprising. In fact, one of the early motivations for NCS [60] was the idea that the termination event in VH-mode penetrates into the core by coupling with low order harmonics that remain large in the core (generally \( \xi \sim r^{m-1} \)). Evidence for this was suggested by the fact that the energy loss is much greater for VH-modes when an \( m/n = 1/1 \) mode is simultaneously present [14,26]. It was speculated [60] that the penetration might be relieved by raising q in the core, since \( m \sim n q \). However, the comparison of the VH-mode and NCS H-mode discharges here, suggests that this does not appear to be the important factor for the peeling-ballooning modes, although raising q does avoid the separate 1/1 mode and eliminates the very large energy losses reported in Ref. [14].
The success of the new model for Type I ELMs, especially in its application to the high performance cases here, rests on the basic assumption of a correspondence between the linear mode penetration and the size of the resulting ELM. The work described here can, in fact, be considered as further verification of the validity of this assumption.

There is little direct data to support this assumption beyond the correlations between the global confinement and $b$ drop and the linear mode width found in a wide range of experimental conditions [22,35-38] and now here as well. This would require fast equilibrium and local confinement data across the entire plasma to use in a reliable transport calculation with all the sinks and sources properly accounted for. Nevertheless, independent experimental evidence now exists for this claim as reported in Ref. [32]. In a study of a VH-mode discharge, the local $T_e$ reduction from Type I ELMs, obtained by averaging over several ELMs, does, in fact, match the envelope of the linear modes computed to be unstable at the time of the first ELM. This result is further encouraging. Future theory can hopefully show why this correlation holds. Superficially, it seems naive. One could reasonably speculate that since the linear mode width is a weak function of $n$, the nonlinear mode also has much the same mode width, which could directly cause the confinement drop. Alternatively, the ELM, comprising the nonlinearly coupled linear modes could trigger another mode or set of modes locally. This may be manifested as an axisymmetric crash. At least in the case when the $m/n = 1/1$ mode is present in the VH-mode discharges, this may be the case, but the crash apparently then penetrates most of the discharge.
V. SUMMARY AND CONCLUSIONS

The termination events from a VH-mode and an NCS H-mode discharge in DIII-D are analyzed in terms of the newly emerging model for Type I ELMs [18–22,30–38] and compared to a conventional H-mode Type I ELM. A suitable relative measure of high performance for these discharges is the peak product $\beta_N H$. For the conventional H-mode discharge this product is about 4. For the VH-mode and NCS H-mode discharges, $\beta_N H \sim 8$ and 6 respectively, but is reduced to ELMing H-mode values after the first large ELM. In this model, the Type I ELM cycle is considered to result from a complex interaction between high $n$ ballooning and low $n$ peeling-balloonning stability, whereby the pressure gradient is limited by either the ballooning limit or an intermediate $n$ peeling-balloonning limit, depending on whether there is access to the second ballooning stability region or not. In turn, that access is controlled partly by the bootstrap current; the bootstrap current increases second stability access but lowers the intermediate $n$ peeling-balloonning mode limits.

For the high performance VH-mode and NCS H-mode discharges, second regime access over a broad region permits a broad pressure pedestal which, in turn, appears to result in a broad, linearly unstable, relatively low $n$, peeling-balloonning instability, that extends well into the half radius. The resulting ELM then destroys the internal high confinement region. This appears to be permanent; the bootstrap current tends to be lost or redistributed and second stability access is then eliminated, preventing the pressure gradient from rebuilding, apparently irreversibly. This may be the result of a change in edge collisionality from a significant flux of impurities into the plasma after the first ELM, in which case, it may be reversible in principle if the influx can be controlled. Cross section shape is also a key determinant of second stability access and this holds further promise of providing some control over the Type I ELMs and their effects. Stronger shaping at least appears to be a key feature of the exceptional high performance discharge ELMs where confinement is not permanently destroyed. Future work will need to investigate in detail the distinctive features of how local confinement and profiles change after the first large ELM in high performance discharges and possibly identify a plausible physical mechanism linking the linear instability mode width and the ELM size.
VI. REFERENCES


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