Physics design point of high-field stellarator reactors

J.A. Alonso*, I. Calvo*, D. Carralero*, J.L. Velasco*, J.M. García-Regaña*, I. Palermo and D. Rapisarda

Laboratorio Nacional de Fusión, CIEMAT, Av. Complutense 40, 28040 Madrid, Spain

E-mail: arturo.alonso@ciemat.es

Received 8 October 2021, revised 20 December 2021
Accepted for publication 10 January 2022
Published 7 February 2022

Abstract

The ongoing development of electromagnets based on high temperature superconductors has led to the conceptual exploration of high-magnetic-field fusion reactors of the tokamak type, operating at on-axis fields above 10 T. In this work we explore the consequences of the potential future availability of high-field three-dimensional electromagnets on the physics design point of a stellarator reactor. We find that, when an increase in the magnetic field strength $B$ is used to maximally reduce the device linear size $R \sim B^{-4/3}$ (with otherwise fixed magnetic geometry), the physics design point is largely independent of the chosen field strength/device size. A similar degree of optimization is to be imposed on the magnetohydrodynamic, transport and fast ion confinement properties of the magnetic configuration of that family of reactor design points. Additionally, we show that the family shares an invariant operation map of fusion power output as a function of the auxiliary power and relative density variation. The effects of magnetic field over-engineering and the $R(B)$ scaling of design points with constant neutron wall loading are also inspected. In this study we use geometric parameters characteristic of the helical axis advanced stellarator reactor, but most results apply to other stellarator configurations.

Keywords: stellarator reactor, stellarator optimization, high-field fusion reactor

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

1. Introduction

In magnetic confinement fusion, the stellarator configuration offers viable solutions for what today appear to be serious problems of the more developed tokamak concept. Since the confining field of a stellarator is created by external coils, without the need to drive large amounts of current within the plasma, continuous operation is granted. This improves the plant availability factor, which positively impacts the overall tritium breeding ratio [1]. Furthermore, since the intrinsic plasma current in a stellarator is typically small, sudden plasma terminations are not expected to limit the reliability or the integrity of a stellarator reactor device. Major disruptions, causing large electromagnetic forces in the structure of the device and generating very energetic beams of run-away electrons, continues to be a concern for tokamak reactors.

The confinement laws of both tokamaks and stellarators display a gyro-Bohm dependence [2]. However, contrary to tokamaks, $H$-mode operation is not considered in baseline stellarator reactor designs. The understanding of $H$-mode transition in stellarators is not any more mature than in tokamaks, but observations suggest that a strong pedestal formation is not typical of $H$-modes in stellarators, with temperature profiles showing no clear edge pedestal. On the one hand, this leads to only mild improvements in confinement for those regimes ($\sim 20\%$) and, on the other, to a more benign edge localised mode activity [3–5]. Operating tokamak reactors in $H$-mode imposes a
minimum power across the separatrix which, in combination with the requirement to radiate a large fraction of the power that enters the scrape-off layer, has been shown to possibly limit the design space of a fusion reactor based on the tokamak concept [6, 7].

The heat exhaust problem, i.e. the need to handle the vast amount of plasma self-heating as it eventually crosses the magnetic separatrix and travels towards the divertor plates, is shared by all approaches to magnetic confinement fusion. The greater diversity of stellarator configurations results in a more diversified scrape-off layer and divertor concepts and geometries. The helical coils of the heliotron device offer a natural long-legged divertor with a well separated volume [8]. The helical axis advanced stellarator (HELIAS) relies on the so-called island divertor, which features multiple X-points and long connection lengths. The operation of an island divertor requires to reduce the plasma current by design, and/or actively control it, to position the field resonance in the plasma periphery. Quasi-axisymmetric and helically-symmetric stellarators feature large bootstrap currents and therefore need to resort to ‘non-resonant’ divertors (see e.g. [9]), that, so far, have not found an experimental realisation. The island divertor concept was first implemented in the Wendelstein 7–AS stellarator [11] and further engineered in its successor, Wendelstein 7-X (W7-X) [12]. First W7-X divertor operation has demonstrated a potential to handle power [13] and impurities [14] favourably. Furthermore, long and stable power detachment has been attained [15, 16]. The characterisation and understanding of the island divertor scrape-off layer physics and plasma wall interaction [17] will occupy the fusion community for years to come, but significant progress is already being made [18–20].

In this and other respects, the stellarator physics basis is under development and the engineering feasibility of a stellarator power plant needs to be demonstrated. The three-dimensional magnetic configuration poses specific challenges for confining the hot plasma fuel and fusion-generated fast particles. It also complicates engineering aspects, which mainly arise from the three-dimensionality of the electromagnets and the need to have enough of them, close enough to the plasma, to generate an optimised magnetic configuration. Superconducting coil manufacture and support structure, breeding blanket geometry or remote maintenance are all more challenging in stellarator reactors [21]. Nevertheless, advances in coil optimization codes are providing new insights and less constraining geometries [22, 23].

In the scientific literature, there exist stellarator reactor studies based on several magnetic configurations, including compact quasi-symmetric [24], force-free heliotron [25, 26] and HELIAS [27–29]. An early comparison of them was published in [30]. Recently, the fusion reactor systems code PROCESS has been adapted to deal will stellarator configurations [31], which will allow to assess economic and technological feasibility of specific 3D equilibria and coil sets.

The physics model used in reactor studies is generally rather coarse-grained, such that only the general geometric parameters of the different configurations (e.g. aspect ratio, rotational transform, size and magnetic field) enter the description through scalings and operational limits inferred from experience in one or several devices. Similarly, for devising design points for stellarator reactors, assumptions need to be made on aspects such as the shape of the density and temperature profiles, the concentration of helium ash or the efficiency of the alpha particle heating. This is simply the reflection of an incomplete understanding of the transport processes in magnetic confinement plasmas. Current stellarator research aims at filling those gaps, understanding the relevant equilibrium, stability and transport physics that would allow identifying credible design points for a stellarator reactor.

The present work builds on these studies, and shares their main methods and assumptions, to assess the impact of a coil technology capable of producing strong confining magnetic fields, stronger than those possible with conventional superconductors like niobium-tin, on the physics design point of a stellarator reactor. The demonstration of the fundamental technological principle, namely, the high temperature superconductors (HTS) of the REBCO type, exist since several decades [32], but has only recently received a broader attention from the tokamak (see e.g. [33–35]) and stellarator [36, 37] fusion communities (see also the recent review [38]). As a first step towards the demonstration of the high-field path to commercial fusion energy, researchers have set off to demonstrate the technological feasibility of high-field HTS magnets for its later use in small-scale proof-of-principles experimental devices [39].

Since both the strength of the magnetic field and device size positively affect the energy confinement time, a trade-off between them is possible. Reference [40] concludes that the ability to produce stronger fields in tokamak configurations (of the order of 10 T on-axis) allows to reduce the size (~5 m major radius), and potentially also the cost, of a fusion reactor while maintaining sufficient confinement for ignition. In doing so, an obvious problem arises: a smaller wall surface leads to larger neutron wall loads and a faster degradation of in-vessel components. This requires not only alternative neutron stopping and tritium breeding approaches, like those based on molten-salt liquid blankets, but also developments in maintenance schemes (e.g. demountable coils) that reduce the reactor down-time [34]. The increase in field and current density on the coils yields larger stresses on the structural components, which can pose important engineering challenges. Furthermore, the unmitigated parallel heat flux towards the divertor is also expected to grow as the magnetic field increases (lowering the power flux perpendicular decay length) and the size is reduced (lowering the linear dimension of the strike lines).

But reducing the linear scale of a stellarator reactor carries additional complications besides those related to higher wall fluxes. The radial extent of the tritium breeding blanket cannot be downscaled, so that the magnetic configuration of smaller devices need to be produced by coils that are proportionally further away from the last closed magnetic surface.

\[ \text{The helically symmetric experiment has been shown to possess a resilient localised pattern of field-line strike points [10], but a divertor was not foreseen in its construction.} \]

\[ \text{The physics model used in reactor studies is generally rather coarse-grained, such that only the general geometric parameters of the different configurations (e.g. aspect ratio, rotational transform, size and magnetic field) enter the description through scalings and operational limits inferred from experience in one or several devices. Similarly, for devising design points for stellarator reactors, assumptions need to be made on aspects such as the shape of the density and temperature profiles, the concentration of helium ash or the efficiency of the alpha particle heating. This is simply the reflection of an incomplete understanding of the transport processes in magnetic confinement plasmas. Current stellarator research aims at filling those gaps, understanding the relevant equilibrium, stability and transport physics that would allow identifying credible design points for a stellarator reactor.} \]
Flux surface shaping and tailoring of the magnetic field spectrum is consequently more challenging and leads to strongly shaped coil designs, since the higher-order modes of the B-field decay quickly when moving away from the coils [22, 41]. One is led to conclude that the exploitation of HTS for fusion applications must be accompanied by important developments in the design and engineering of a stellarator reactor. These developments could certainly build on the above-mentioned proof-of-principle solutions in the tokamak line. We note that stellarator reactors based on catalysed D–D fusion and low-temperature superconductors have also been studied [42].

In the present work we address the questions: how would very high field HTS-based electromagnets impact the physics design point of a deuterium-tritium stellarator reactor? What stellarator physics research directions could prepare us for making use of such a technological development? We make note that it is a premise of this study that the many engineering challenges associated with the fabrication of strongly shaped electromagnets and their operation at very-high fields are surmounted. Although we discuss some engineering aspects of higher-field, smaller devices, the focus of this study lies on the physics characteristics of the reactor design point.

The rest of the paper is organised as follows. In section 2, we present the device geometry, physics assumptions, and formulas that will be used in the one-dimensional study of stellarator reactor design points. A simplified zero-dimensional analysis is first conducted in section 3, where we derive the basic field strength/device size relation that stems from the empirical scaling of energy confinement time. We show that this relation leads to the approximate invariance of several important physical parameters for a family of reactor design points. Importantly, the scaling of the density operation point is introduced in this section (and further elaborated in Appendix B). In section 4, we complement the findings of the previous section with 1D profile analysis based on prescribed profile shapes. The 0D invariances translate into several archetypal plasma profiles that are shown and discussed. The effect of magnetic field over-engineering on the physics design point is also discussed in this section. Our main conclusions are summarised in section 5.

2. Identification of stellarator reactor design points in the \((B, R)\) plane

In this section we present the basic parameters, formulas and hypotheses that are used to identify potential reactor design points in the \((B, R)\) plane. Here \(B\) is the characteristic on-axis magnetic field strength and \(R\) is the major radius of the device. We will be working with a fixed device aspect ratio and magnetic geometry, so that \(R\) characterises the device size and determines all other dimensions such as the minor radius of the plasma column, \(a\), or the confinement volume \(V_a = 2\pi^2 Ra^2\).

Given certain values for \(B\) and \(R\), and other confinement-relevant configuration parameters, we would like to find plasma parameters that fulfil the simplified version of the power balance equation,

\[ P_h = \frac{W}{\tau_E} \]  \(\text{(1)}\)

where \(P_h\) is the net heating power, \(\tau_E\) the energy confinement time and \(W\) the total internal plasma energy given by the volume integral

\[ W = \frac{3}{2} \sum_s \int d^3 x n_s T_s, \]  \(\text{(2)}\)

where the integration domain is the volume within the last closed magnetic surface. The index \(s\) labels the plasma species with particle density \(n_s\) and temperature \(T_s\). In the rest of the paper, we will assume that all species have the same temperature and therefore drop the species index and refer to the plasma temperature \(T\). Plasma is composed of electrons, with number density \(n_e\), deuterium and tritium ions and helium ash from the DT fusion reactions. We assume equal amounts of D and T and a 5% He concentration, i.e., \(n_{He} = 0.05n_e\), so that \(n_D = n_T = 0.45n_e\).

The energy confinement time, \(\tau_E\) in equation (1) is assumed to be given by the ISS04 scaling [43],

\[ \tau_E = f_c \times 0.134 \times 28^{0.64} R^{0.61} n_e^{0.54} B^{0.84} T^{0.41}, \]  \(\text{(3)}\)

where \(f_c\) is the configuration factor, \(n_e\) is the line-averaged electron density and \(T_{2/3}\) is the value of the rotational transform at the \(\rho = 2/3\) magnetic surface. The normalised radius \(\rho\) as well as the radius \(r\) will be used to label magnetic surfaces in this work, which relate to each other and to the volume within the flux surface, \(V\), as \(\rho = r/a = \sqrt{V/V_\alpha}\). In terms of these variables, a volume integral, like the one in (2), is given by

\[ \int_0^{V_\alpha} dV f(V) = 2V_\alpha \int_0^1 d\rho \rho f(\rho) = \frac{2V_\alpha}{a^2} \int_0^a dr f(r), \]

whereas the line-averaged electron density in (3) is given

\[ n_e = \int_0^\rho n_e(\rho) d\rho. \]

Note that all physical magnitudes are assumed constant on flux surfaces in this study.

For estimating the net heating power, \(P_h\), we consider the plasma self-heating by alpha particles, the power lost by Bremsstrahlung radiation and any auxiliary external heating power. Each of these terms is modelled with a \(\rho\)-dependent power density, \(S_a, S_B\) and \(S_{aux}\) that are given below:

\[ S_a = E_a n_D n_T (\sigma v)_{DT}(T), \]  \(\text{(4)}\)

with \(E_a = 3.5\ \text{MeV}\). The form of the DT reactivity, \((\sigma v)_{DT}\), can be seen in figure 1. It displays an approximate quadratic temperature dependence in the range of interest. The local Bremsstrahlung radiation density is given by

\[ S_B = C_B Z_{eff} n_e^2 \sqrt{T}. \]  \(\text{(5)}\)

If temperature and density are expressed in keV and \(10^{20}\ \text{m}^{-3}\) respectively, the numerical constant \(C_B = 5.35 \times 10^{-3}\) gives

\(^2\) We will later show that this concentration of helium implies a ratio of He particle to energy confinement time of about 9 for the family of reactors discussed in section 4.
the Bremsstrahlung radiation density in units of MW m\(^{-3}\). The effective ion charge \(Z_{\text{eff}}\) is defined as \(Z_{\text{eff}} = \sum_{k} Z_{k} \rho_{k} / \sum_{k} \rho_{k}\), and is equal to 1.1 with our assumed plasma composition. The auxiliary power is modelled by a centred Gaussian

\[
S_{\text{aux}} = \frac{P_{\text{aux}}}{V_{\text{a}} w} e^{-\frac{\rho^2}{w^2}},
\]

where \(P_{\text{aux}}\) is the total auxiliary power and \(w\) is the width of the power deposition profile. The volume integrals of the power densities (4) and (5) are termed \(P_{\alpha}\) and \(P_{\beta}\) respectively.

The net heating power is then defined here as \(P_{\text{h}} = \max(P_{\alpha})\), where

\[
P(\rho) = 2 V_{\text{a}} \int_{0}^{\rho} d\rho \rho (k_{\alpha} S_{\alpha}(\rho) + S_{\text{aux}}(\rho) - S_{\beta}(\rho)).
\]

The \(k_{\alpha}\) constant is the efficiency of the alpha heating. We take \(k_{\alpha} = 0.9\), thereby allowing for a 10% loss alpha particle energy due to fast ion transport\(^3\).

For the purpose of this study, a reactor design point will be a point in the \((B, R)\) plane for which a device producing a fusion power \(P_{\text{ fus}} = 3\) GW with a fusion gain \(Q = P_{\text{ fus}} / P_{\text{ aux}} = 40^4\)

\(^3\)Lower orbit losses are achievable through optimization in time-independent fields (see e.g. the recent study in [46]), but turbulence and/or Alfvén modes can enhance fast particle transport. The specific choice of \(k_{\alpha}\) is not important for the conclusions of this work, as far as it stays sufficiently close to 1.

\(^4\)Again, this arbitrary choice does not affect our conclusions. The reactor design point does not depend strongly on the desired \(Q\), for large enough \(Q\) values (see figure 2(j)). Lowering the required auxiliary power would improve the balance of plant but could make active burn control by density necessary. Similarly, the lower recirculating power of stellarator power plants could justify targeting lower \(P_{\text{ fus}}\) values (see e.g. [47]). We choose to stay with 3 GW to allow a more straightforward comparison with other reactor design points in the stellarator literature.

**3. Simplified 0D analysis: \(B(R)\) scaling of a reactor design point and its main plasma parameters**

Before presenting the analysis of reactor design points based on the formulas introduced in the previous section, we carry out next a simplified 0D analysis that will guide intuition and help understand those results. We will identify the \(B(R)\) dependence of a reactor design point and the ensuing dependency of the main physics dimensionless parameters on \(B\) (or \(R\)).

**Table 1. Reactor design parameters kept fixed in this study. Note that the fusion power accounts for alpha and neutron energy, as well as for exothermic breeder reactions [48, 49].**

| Parameter name            | Notation | Value |
|---------------------------|----------|-------|
| Aspect ratio              | \(A\)    | 9.0   |
| Rotational transform \((\mu = 2/3)\) | \(\mu/3\) | 0.9   |
| Configuration factor      | \(f_{\text{c}}\) | 1.0   |
| Helium concentration      | \(n_{H}/n_{a}\) | 0.05  |
| Fuel concentration        | \(n_{D}/n_{a}\) | 0.45  |
| Effective charge          | \(Z_{\text{eff}}\) | 1.1   |
| \(\alpha\)-heating efficiency | \(k_{\alpha}\) | 0.9   |
| Fusion-to-alpha power ratio | \(P_{\text{ fus}}/P_{\alpha}\) | 6.4   |

**Figure 1. Reactivity of the D–T fusion reaction as a function of the fuel temperature. The points (+) are calculated by convolving the Maxwellian distributions with the cross section \(\sigma\) given by Duane’s parametrization [44]. The blue curve is the parametric fit to the reactivity given by Bosch and Hale [45]. The red parabola approximates the reaction fairly well below 20 keV (the multiplying coefficient is chosen to match the reactivity at 10 keV).**
Figure 2. Magnetic field $B$ and device major radius $R$ scan of several physics and engineering parameters. The line-averaged electron density is initially set to the value given by (11) and the auxiliary power is determined to give a total fusion power of 3 GW, shown in subfigure (g). For small, low-field devices this condition cannot be met (points marked with 0 GW in (g)). For the $Q = \infty$ part of the plane (where $P_{\text{aux}} = 0$) the density is lowered to meet the 3 GW target. Marked in red is the $Q = 40$ line that can be considered for design points of stellarator reactors. We note that the 1D study conducted later in the paper will lower this curve by about 15%.

In order to do that, we will introduce an ad hoc scaling of the plasma density that will also be used later in this article.

For the 0D analysis of this section, we rewrite equation (2) using $\sum n_i = 2 n_e$ and defining $\langle n_e T \rangle_V = \frac{1}{V} \int_V n_e T \, dV$, to get

$$\langle n_e T \rangle_V \approx \frac{P_{\text{th}} V_e}{3 V_d}. \tag{8}$$

A 0D plasma beta is then defined as

$$\beta_0 = \frac{2\langle n_e T \rangle_V}{B^2/2\mu_0}. \tag{9}$$
whereas a characteristic temperature is obtained from the relation
\[ T_0 = \frac{\langle n_e T \rangle}{\pi \rho_0}, \tag{10} \]
which is assumed to be close to the volume-averaged temperature. This temperature is used to estimate the alpha and the Bremsstrahlung power. In equation (8) we include the alpha and auxiliary heating in \( P_h \), but otherwise neglect Bremsstrahlung radiation losses.5

To proceed with the analysis one needs to determine how line-averaged electron density, \( \pi_e \), scales with \( R \) and \( B \). In this work we choose density to scale like
\[ \pi_e (m^{-3}) = 1.04 \times 10^{19} (B(T))^2, \tag{11} \]
where the pre-factor is chosen to match the HSR4/18 density (2.6 × 10^20 m^{-3} at 5 T). This choice, which fundamentally determines the conclusions of this article, is consistent with the increase of the cut-off densities of the electron–cyclotron resonance heating an it results in a constant ratio of density to the critical density, \( \pi_e / \pi_{ec} \), along the lines constant fusion power and gain in the (B,R) plane. Here \( \pi_{ec} \) is the line-averaged critical electron density determined by a Sudo-like radiative limit. Further elaboration of this choice is given in appendix B.

The result of the 0D (B,R) scan is shown in figure 2. The plots in this figure are made as follows: with the density initially set to (11) we proceed, for each \( (B,R) \) pair, to find the plasma temperature required to meet the \( P_{\text{ fus}} = 3 \) GW target (figure 2(g)) together with the auxiliary power (figure 2(h)) that is needed to get to that temperature, according to the \( \tau_E \) scaling (3). For those points for which this process leads to \( P_{\text{ aux}} < 0 \), the density is lowered until \( P_{\text{ fus}} = 3 \) GW and \( P_{\text{ aux}} = 0 \) (this corresponds to the \( Q = \infty \) part of the plane). The resulting \( (n_e, T) \) pairs are shown in the plots figures 2(a) and (b). The rest of the physical and engineering parameters (figures 2(c)–(l)) are derived from them and the \( (B,R) \) values. The definitions of the normalised tritium gyro-radius and collisionality (figures 2(d) and (e)) are standard and are presented in appendix A.3 for convenience. The consequences of increasing \( B \) at a fixed device size while staying with a 3 GW fusion target can be seen in figure 2: the fusion gain increases rapidly reaching \( Q = \infty \). Beyond this point the density operation point needs to be lowered to moderate the increase in the confinement time and reduce the fusion reaction rate. Temperature needs to be increased, which reduces the collisionality. The plasma beta is lowered by the increase in the field strength. While the second of these reductions alleviates the MHD equilibrium and stability requirements on the magnetic configuration, the reduction of density and collisionality makes it more difficult to confine fast and thermal particles in the three dimensional configuration. In consequence we conclude that, unless the reduction of \( \beta_0 \) design point is a strong requirement, the optimum location of design points for the stellarator reactor in the \( B,R \) plane are those just below the \( Q = \infty \) line. The red line shown in all plots in figure 2 are the \( Q = 40(B,R) \) pairs, which are considered potential reactor design points in this work.6

It is apparent in this figure that points along constant-\( Q \) lines (or \( Q \)-lines) feature a very similar temperature \( T_0 \) (figure 2(b)). This is a consequence of the \( n \sim B^2 \) scaling, for the 0D temperature (10) scales as \( T_0 \sim \tau_E / n V_a (P_h \) held constant). According to (3), this gives a weak
\[ T_0 \sim (RB)^{-0.08} \tag{12} \]
dependence. Together with a constant fusion power, \( P_{\text{ fus}} \propto n^2 V_a (\sigma V)_{DT} \), this results in the \( B(R) \) relation,
\[ B \sim R^{-3/4}, \quad (\text{constant } P_{\text{ fus}} \text{ and } Q), \tag{13} \]
where the weak dependence (12) has been neglected. Note that the aspect ratio is assumed constant throughout this article, so that \( V_a \sim R^3 \). The \( Q = 40 \) line in the plots in figure 2 closely follows this \( B(R) \) dependence.

The scaling (13) can be understood as the maximal reduction in device size possible by an increase in the magnetic field strength. That is, given a reactor design point \( (R_0, B_0) \), the ability to generate a larger field \( B_1 \) allows to reduce the relative reactor size as \( R_1 / R_0 = (B_0/B_1)^{4/3} \). Further reducing the device size would require to increase the \( \pi_e / \pi_{ec} \) ratio to recover the \( P_{\text{ fus}} = 3 \) GW, \( Q = 40 \) target. It can be shown that the scalings (12) and (13) stem from the approximate gyro-Bohm dependence of the ISS04 energy confinement time (3), together with the \( n \sim B^2 \) scaling used here. In fact, it is straightforward to check that balancing a gyro-Bohm diffusive heat flux, \( \Gamma_Q \propto n \chi_{gb} d T / d r \) (see appendix A.2), against a heating power divided by the flux-surface area \( \sim P_h / R^2 \) yields a constant temperature, independent of \( B \) and \( R \), when the heating power and the normalised temperature scale length are kept constant. The \( B(R) \) scaling (13) then follows from imposing a constant alpha power with a fixed temperature.

Several other important physics parameters are kept almost constant along the \( Q \)-lines in figure 2. The Bremsstrahlung power has a dependence similar to that of the alpha power, \( P_B \sim n^2 V_a f(T_0) \), where \( f(T_0) \sim \sqrt{T_0} \), and is similarly constant along the \( Q \)-lines. Plasma beta, \( \beta_0 \sim n T_B / B^2 \), is approximately constant along the \( Q \)-lines, whereas triton collisionality \( (\nu_T^* \sim n R / T^2 \sim R^{-1/2}) \) and gyroradius \( (\rho_T^* \sim 1 / RB \sim R^{-1/4}) \) vary only slowly with \( R \) (or \( B \)). The combination \( (\nu_T^*)^{-1}(\rho_T^*)^2 \) is approximately constant along the \( Q \)-lines. On the contrary, at least two engineering parameters, the neutron wall loading (NWL) and the magnetic energy (figures 2(k) and (l)), do show dependencies on the device size. The NWL decreases inversely to the first wall surface area \( \sim R^2 \) when the neutron power is kept constant, as in this scan. The values shown in the first of these figures have been calculated assuming an average plasma-wall distance required for hosting the diverting

5 In a 0D treatment like the one used in this section, the use of an average temperature like (10) leads to an overestimation of \( P_h / P_a \).

6 We note that the 1D study conducted later in the paper will shift the \( Q = 40 \) line towards lower fields/smaller device size. In section 4.5 we will further discuss the consequences of stepping out of this line.

7 Since all densities are related to each other by constant proportionality factors in this work, we will sometimes use \( n \), without subindex, when discussing scalings with either electron or ion densities.
magnetic structure equal to 20% of the plasma minor radius \( a \). We defer a longer discussion about the implications of these neutron fluxes and other engineering aspects of high-field devices to section 4.4. The vacuum magnetic energy is estimated by \( E_B = V_cB^2/2\mu_0 \), where the volume enclosed by the coil set of average radii \( c \) is approximated by \( V_c = 2\pi^2Rc^2 \). \( E_B \) decreases for smaller, higher-field reactors somewhat slower than \( \sim R^3B^2 \sim R^{3/2} \) along the \( Q = 40 \) line, since the thickness of the neutron shield and beryllium blanket is kept equal to 1.3 m in the scan\(^8\).

4. 1D analysis of reactor design points with prescribed \((n, T)\) profiles

The 0D analysis presented in the previous section allowed to identify dependencies of the main design parameters that are implied by the ISS04 energy confinement time. We found that a quadratic scaling of the density, \( n \sim B^2 \) leads to a relation \( B \sim R^{-3/4} \) of the size and field strength of stellarator reactors with the same fusion energy and gain. This, in turn, leads to the invariance of several physics parameters for a family of reactor design points. However, the precise identification of a design point requires considering profile effects, for most of the fusion reactions are produced in the hottest central part of the plasma column. In this section we will carry out a 1D study of reactor design points and show how the invariances derived in the 0D analysis lead to archetypal profiles for several plasma parameters. Furthermore, we will show that the reactor operation map, \( P_{\text{aux}}(P_{\text{aux}}, n/n_{\text{DP}}) \), is also shared by the family of stellarator reactors. Here \( n_{\text{DP}} \) is the density at the design point.

The formulas that will be employed for the 1D analysis were presented in section 2, and involve plasma density and temperature profiles. To date, no validated first-principle transport model has been developed that can be used to predict these profiles in a stellarator reactor confidently. Neoclassical transport theory and codes are well established and have been tested, often positively, in 3D magnetic configurations. Neoclassical heat fluxes provide an irreducible transport, found to be generally smaller than the experimental fluxes \([50, 51]\). Turbulence is thought to cause the additional heat transport, the computation and validation of which is an active subject of current research \([52–54]\). An empirical parametrization of turbulent fluxes based on experimental data was used in \([29]\). The analysis presented here will be conducted with prescribed shapes of density and temperature profiles. The line-averaged density is given by \( \langle n \rangle \), whereas the temperature profile is scaled to provide consistency with the power balance \( (1) \) and the ISS04 energy confinement time \( (3) \). Details on the form of these profiles are given in appendix A.1. We note that the shape of the temperature profile is chosen to depend on the particle and heating power densities with a constant-\( \chi \) ansatz.

The effect of varying the density profile shape on the reactor design point will be briefly inspected in section 4.3. An example of the \( n \) and \( T \) profiles is shown in figures 3(a) and (b), which approximates the design point of the HSR4/18 device with \( R = 18 \) m, \( B = 5 \) T and \( P_{\text{aux}} = 3 \) GW, \( Q = 40 \). Other plasma parameters, like the volume-averaged plasma beta \( \langle \beta \rangle \), the energy confinement time \( \tau_E \), or the total power radiated through Bremsstrahlung \( P_B \), closely resemble design values reported in \([28]\). Note that the scaling of the density \( (11) \) is indeed chosen to match the HRS4/18 line-averaged density at its design field of 5 T.

The resulting alpha heating and Bremsstrahlung power density profiles are shown in figure 3(c), together with the auxiliary heating density of width \( w = 0.36 \) (see equation \( (6) \)). Figure 3(d) shows the plasma and fast particle beta, \( \beta_\alpha \), the tritium burnup profile and total burnup fraction, \( f_{\text{burnup}} \). The classical slowing-down distribution function of alpha particles is used to obtain an estimate for their pressure and beta (see appendix A.4), which is important to determine the properties of the Alfvén spectrum and the Alfvén-induced energetic particle transport (see e.g. \([55]\) for a review). Recently, a potential stabilising effect of the fast particle pressure on the ion temperature-gradient driven turbulence has been put forward in \([56]\).

The tritium burnup fraction is defined as the number of DT reactions in the plasma per second over the number of injected tritium atoms per second (see appendix A.5). The inverse of the burnup fraction can be understood as the average number of times that a tritium atom needs to be cycled through the vacuum vessel before it undergoes a DT reaction. Each of the times that a tritium atom is cycled, there is a certain probability that it be lost from the fuel cycle. The burnup fraction is therefore an important factor to determine the overall tritium breeding ratio. It is, nevertheless, subject to considerable uncertainty. Nishikawa sets a \( f_{\text{burnup}} > 0.5\% \) requirement to enable tritium self-sufficiency \([57, 58]\), upon consideration of production of tritium in the blanket system and several losses due to trapping, permeation and decay in the vacuum vessel and fuelling and storage circuits. The same references quote values around 3%–4% to ease the requirements on the efficiency of tritium recovery and breeding. A recent review \([1]\) concludes that tritium self-sufficiency can be achieved with confidence with \( f_{\text{burnup}} > 2\% \) together with a plant availability factor greater than 50%. In figure 3(d) we estimate the burnup fraction using several, partially compensating approximations; namely, zero recycling, 100% fuelling efficiency and an equal tritium particle and energy confinement times. This results in an estimated \( f_{\text{burnup}} = 1.2\% \), which falls in the right range\(^9\).

Finally, fundamental quantities related to the neoclassical and turbulent transport are plotted in figure 3(e) (normalised tritium gyro-radius and collisionality profiles) and figure 3(f).

---

\(^{8}\) It should be noted that the simplistic estimation of the average coil radius, \( c = 1.3 \ m + 1.2a \), results in a considerably smaller \( c \) and total magnetic energy \( E_B \) than those quoted in \([28]\) for \( a = 2 \) m. Missing a specific MHD equilibrium and coil design, the values and tendencies shown in figure 2(l) should be considered indicative.

\(^{9}\) Since there is a large uncertainty in this estimate, the quantitative burnup fractions presented in this work should be de-emphasised. More important for the discussion is the fact that, under our assumptions, this fraction does not depend on the reactor size (see below).
Figure 3. Plasma profiles for an $R = 18$ m, $P_{\text{in}} = 3$ GW, $Q = 40$ stellarator reactor. The average values obtained are very close to those of the HSR4/18 (cf [28]).

(temperature and density scale lengths and the thermal diffusivity at mid radius, $\chi_{\alpha_5}$; see appendix A.2).

The $(B, R)$ position of the design point with profiles shown in figure 3, is plotted in figure 4 together with those of smaller, higher-field devices and a larger $R = 22$ m device. These are calculated similarly, keeping the auxiliary power constant at $P_{\text{aux}} = 75$ MW and varying the magnetic field $B$ until the fusion power target, $P_{\text{in}} = 3$ GW, is met. As found in the previous section, the resulting design points fall very close to the $B \sim R^{-3/4}$ curve. Table 2 summarises the characteristics of those design points. We note that several global magnitudes not listed in that table are nearly constant for the five design points shown in the table (as well as for the other points along the $Q = 40$ line in figure 4) as anticipated in section 3. This includes the volume averaged plasma $\beta$, the net heating $P_\beta$ and Bremsstrahlung power $P_B$, the fusion burnup fraction $f_{\text{burnup}}$ and the ratio of helium particle confinement time $\tau_{\text{He}}$ to the energy confinement time (see caption in table 2). Furthermore, we will show next that the constancy holds as well for the profiles of certain physics parameters.

At this point we note that, as already stated in [29], the synchrotron radiation is not a relevant loss of power for the usual HELIAR reactors design points. This is also the case for the higher field points considered here. In fact the vacuum synchrotron emission scales as $\sim V_e B^2 n_e T_e$, which is constant along the $Q = 40$ line in figure 4. The plasma opacity factor depends on the density, field and plasma size as $\sim a n_e / B$ (see e.g. [59]), which slightly increases with the size of the reactor, $\sim R^{1/4}$, along that line. Also note that, in general, the stellarator design points feature higher densities and lower temperatures compared to the tokamak’s, which result in substantially lower synchrotron radiation losses for a similar volume and field strength.

4.1. Invariant profiles and common optimization requirements

There are a number of important plasma profiles that are approximately the same for all the reactor design points with $P_{\text{in}} = 3$ GW and $Q = 40$ (dashed line in figure 4). This was anticipated in the 0D analysis of section 3 and is recovered here using a 1D analysis (see figure 5). Shown in this figure are profiles for selected plasma parameters computed for the five design points listed in table 2. Line thicknesses show the maximum differences in the profiles obtained from the 1D power balance described in section 4. These can be considered archetypal profiles for the reactor family that includes the HSR4/18. It should be noted that details of these profiles do depend on specific choices, like the model density and temperature profiles. Their constancy, however, depends solely on the assumption that the same profile shape can be used to represent density for all reactor sizes and on the ISS04 and $n \sim B^2$ scalings. The temperature profile (figure 5(a)) is found to be almost constant, and so is the temperature scale-length $a / L_T = T^{-1} d T / d \rho$ ($a / L_n$ is also constant by assumption). Both the thermal and alpha-particle $\beta$ profiles (figure 5(b)) and the diffusivity in gyro-Bohm units (figure 5(c)) are nearly invariant. Figure 5(d) shows the combination $\rho T^2 / v_T^2$ of normalised gyro-radius $\rho_T$ and collisionality $v_T$. The fact that these profiles are independent of the reactor size (or field) as indicated above, leads to the
4.1.1. MHD optimization. That similar equilibrium currents, $\beta$ limits, and MHD stability properties are required is implied by the invariance of the $\beta(\rho)$ profile in figure 5. Modifications to the vacuum magnetic field produced by the plasma current densities $j$ are given by the equation

$$\nabla \times \delta B = \mu_0 j.$$

The equilibrium plasma currents include the diamagnetic, parallel Pfirsch–Schlüter and parallel bootstrap current, which can be written as

$$j = \frac{dp}{d\rho} \left( \frac{B \times \nabla \rho}{|B|^2} + hB \right) + \left( \frac{j \cdot B}{|B|^2} \right) B,$$

where $h$ is a function of the space and the magnitude of the magnetic field, defined such that it cancels the total divergence of the term in brackets. This first term, containing the diamagnetic and Pfirsch–Schlüter currents, scales as $\sim (d\beta/d\rho)B/R$ which, plugged into (14), leads to $\delta B/B \sim (d\beta/d\rho)$. In the absence of external current drive, the second term in (15) is given by the bootstrap current. While an explicit expression for the bootstrap current, valid for any collisionality regime, does not exist, its asymptotic form in the $1/\nu$ regime [60] displays a scaling $\langle j \cdot B \rangle \sim nT/R$, which, similarly leads to a relative field modification, $\delta B/B \sim \beta$, independent of the device size. A similar scaling follows from an earlier work by Boozer [61].

4.1.2. Transport optimization. The transport properties of the devices that share the invariant profiles in figure 5 should also bear important similarities. For all of them, the total thermal diffusivity, $\chi$, must be a similarly small fraction ($\sim 10^{-2}$) of the gyro-Bohm diffusivity, $\chi_{GB}$, in the central part of the plasma, which will be shown to be consistent with the expected dependencies in the neoclassical and turbulent diffusivities.

Note that this order of magnitude is implied by the need to reach a sufficient core temperature to achieve the target fusion power of 3 GW and does not depend strongly on the profiles chosen in this study.

Neoclassical energy transport is deemed an important, possibly dominant, transport mechanism in high-temperature stellarator reactor cores. In 3D fields, it features a specific regime of high diffusivity that is inversely proportional the collisionality. In this so-called $1/\nu$ regime, the heat diffusivity of a plasma species $s$ is $\chi_1/\nu \propto (a^2 v_\text{th}/R)(\epsilon_{\text{eff}}/\nu_s^2/\nu_s^2) = (a/R)(\epsilon_{\text{eff}}/\nu_s^2)\chi_{GB}$, where the proportionality constant is of order unity (see e.g. [62]). The effective ripple, $\epsilon_{\text{eff}}$, is a characteristic of the magnetic field structure, the minimisation of which is targeted by stellarator optimization. An order-of-magnitude figure for how small this coefficient needs to be, can be estimated by imposing $\chi \geq \chi_1/\nu$ for the tritium ions. Figure 5(c) shows that $\chi/\chi_{GB} \sim 3 \times 10^{-2}$ is necessary in the centre of the family of stellarator reactors. Together with the weakly size-dependent value of the central tritium collisionality in table 2 this conditions leads to a central $\epsilon_{\text{eff}} \lesssim 1\%$. Values around or below this are characteristic of several existing stellarator configurations (see e.g. [63] and references therein), including the HSR4/18 [28], new compact quasi-axisymmetric [64] and several W7-X configurations. In fact, the analysis conducted in [29] shows that neoclassical transport is compatible with fusion conditions for a scaled version of the high-mirror W7-X configuration. It should also be noted that the radial electric field moderates ion neoclassical losses at the core collisionality, where other regimes like the $\nu$ become increasingly important. This makes the $1/\nu$ estimates a worst-case scenario for the neoclassical transport channel.

The approximate invariance of scale lengths $a/L_T$, $a/L_n$ and $cT^{-1}(d\phi/d\rho)$, an assumed constant magnetic geometry and the only weak size-dependence of collisionality ($\sim R^{-1/2}$, see the end of section 3) and $\nu_s^2 \sim R^{-1/4}$ leads to the conclusion that $\chi_{neo}/\chi_{GB}$ would be similar for all reactor design points. Nevertheless, it is well known that microturbulence enhances energy transport above neoclassical levels. In neoclassically optimized stellarators like W7-X, the turbulent component can even dominate the radial energy fluxes over the entire plasma volume [51]. Given a magnetic configuration,
\( \chi_{\text{b}}/\chi_{\text{GB}} \) depends on \( a/L_T, a/L_m \), and the collisionality. In the family of stellarator reactors discussed in this paper, \( a/L_T, a/L_m \) do not vary, whereas the collisionality varies only slightly. One is then similarly led to conclude that, \( \chi_{\text{b}}/\chi_{\text{GB}} \) would not vary significantly within this family.\(^{10}\)

To conclude the discussion on transport, we note that the fact that \( \chi_{\text{neoc}}/\chi_{\text{GB}} \) and \( \chi_{\text{b}}/\chi_{\text{GB}} \) do not vary much within the reactor family is consistent with the approximate invariance of the required \( \chi/\chi_{\text{GB}} \) (figure 5(c)). This does not come as a surprise, for the ISS04 energy confinement time displays an approximate gyro-Bohm dependence that is also characteristic of both neoclassical and turbulent transport mechanisms.\(^{11}\)

4.1.3. Fast ion optimization. Collisionless fast ion losses in stellarator proceed in two time-scales. Trapped particles cause prompt losses, which have a very short characteristic time given by \( v_{\text{M}}/a \), where \( v_{\text{M}} \) is the characteristic size of the radial magnetic drift. Since this is much shorter than the collisional slowing-down time, the magnetic configuration needs to be designed to reduce prompt losses for any of the reactor sizes considered in figure 4. On a longer time-scale, fast ion losses due to stochastic diffusion set in, which can be moderated by lowering the collisional slowing-down time. A diffusion coefficient for the stochastic losses was derived in [66] using an analytical representation of the magnetic field spectrum. Aside from size-independent geometric factor, the stochastic diffusion scales as \( D_{\alpha} \sim R^2 \Omega_{\alpha}(\rho_{\alpha})^2 \), where \( \Omega_{\alpha} \) is the alpha particle cyclotron frequency. In terms of our size \( R \) and field strength \( B \) variables, one gets \( D_{\alpha} \sim R^{-2} B^{-3} \), which reflects a reduction in the diffusivity of alpha particles for larger devices and stronger fields. To compare the importance of stochastic losses for the different design points, we use an estimate of the diffusive radial excursion in a slowing-down time, \( \langle \Delta \rho \rangle \sim a^{-1} \sqrt{D_{\alpha} \tau_\alpha} \). For constant electron temperature, the slowing-down time \( \tau_\alpha \) is inversely proportional to the density (see equation (A.8)). Using the density and \( B(R) \) scaling (equations (11) and (13)) leads to \( \langle \Delta \rho \rangle \sim R^{-1/8} \). By this measure, stochastic diffusion of alpha particles is expected to be of very similar importance for the various reactor design points considered in figure 4 and table 2.

Alpha particles can excite Alfvén modes which, in turn, can enhance their radial transport. The characteristics of this interaction and enhanced transport would be similar for the different reactor design points in the following sense: first, the population of alpha particles displays an invariant \( \beta_{\alpha} \) profile shown in figure 5 (see also appendix A.4). Second, the Alfvén speed \( v_{\alpha} = B/\sqrt{\rho_{\alpha} \sum_s m_s} \), where \( m_s \) is the mass of the species \( s \), is constant according to the density scaling (11).

---

10 The neoclassical and turbulent diffusivities \( \chi_{\text{neoc}} \) and \( \chi_{\text{b}} \) are defined as the coefficients that relate the size of the corresponding energy flux with the typical scale length of the temperature profile. Namely, \( Q_{\text{neoc}} \sim \chi_{\text{neoc}} n T L_T^{-1} \) and \( Q_{\text{b}} \sim \chi_{\text{b}} n T L_T^{-1} \).

11 This is not to say that deviations in particular parametric dependencies of the neoclassical and turbulence diffusivities with respect to the gyro-Bohm diffusivity are excluded. The neoclassical 1/\( v_T \) limit discussed before is an example of this. Collisional stabilization of the trapped electron codas has been shown to possibly lead to an isotopic mass dependence inverse to that of gyro-Bohm [65]. What is important from the previous discussion is that the other parameters that define the neoclassical and turbulent transport regimes, and might introduce deviations with respect to the gyro-Bohm diffusivity, are themselves nearly equal across the reactor family..
fusion power dependence on the auxiliary power is important characteristic; namely, \( f_i \) in the low \( f_i \) field. The steep variation of fusion power equilibrium temperature profile is greatly increased for the high auxiliary power and density allow to recover the \( P_{\text{aux}} = 3 \) GW; \( Q = 40 \) operation point. However, the operation landscape is substantially changed. The sensitivity of \( P_{\text{aux}} \) is increased for the peaked density profile, such that the operation point lies in a thermally unstable region (in the sense described previously) and sits close to the jump to a higher temperature solution of the power balance equation. It needs to be noted that the operation map is calculated assuming that the D–T fuel mix can be kept constant. In practice, the accumulation of He ash would likely lead to the moderation of the reaction rate for the higher fusion powers.

4.3. Effect of the profile shapes

The shape of the density profile has been so far fixed to that shown in figure 5, with a mild peaking given by \( k_2 = 0.5 \) in equation (A.1). The shape of the temperature profile (A.3) also depends on this choice. In a reactor, the density profile will be determined by the fuelling method and the particle transport characteristics. Peaked density profiles are thought to be beneficial for confinement in W7-X [51]. In the absence of a core particle source, neoclassical thermal diffusion is however expected to lead to core particle depletion [67]. Hollow density profiles have not been reported in W7-X so far, but are common in the LH D device.

For inspecting the effect of the profile shape on the operation point of a reactor we reproduce the operation map of figure 6 for a hollow \((k_2 = -1)\) and a more peaked \((k_2 = 5)\) density profiles. This is shown in figure 7, which shows that relatively small adjustments of the auxiliary power and density allow to recover the \( P_{\text{aux}} = 3 \) GW; \( Q = 40 \) operation point. However, the operation landscape is substantially changed. The sensitivity of \( P_{\text{aux}} \) is greatly increased for the peaked density profile, such that the operation point lies in a thermally unstable region (in the sense described previously) and sits close to the jump to a higher temperature solution of the power balance equation. It needs to be noted that the operation map is calculated assuming that the D–T fuel mix can be kept constant. In practice, the accumulation of He ash would likely lead to the moderation of the reaction rate for the higher fusion powers.

![Figure 6](image.png)

**Figure 6.** Operation map \((P_{\text{aux}}, n/\text{DP})\) for the HSR/4/18 reactor family (dashed line in figure 4). Thermal stability refers to the sign of \( P_h - W/\tau_E \) caused by a \( 1\% \) constant variation of the equilibrium temperature profile. The steep variation of fusion power in the low \( P_{\text{aux}}, \) high \( n/\text{DP} \) region is due to the appearance of a higher temperature solution of the \( P_h = W/\tau_E \) power balance. Note that the fuel mix is kept constant in this scan.

4.2. Invariant operational map

The family of reactors in figure 4 and table 2 share another important characteristic; namely, the operation map (i.e. the fusion power dependence on the auxiliary power \( P_{\text{aux}} \) and relative density variation with respect to the design point \( n/\text{DP} \)), shown in figure 6. That is this indeed the case shown can be from the dependences \( P_{\text{aux}} \sim V_m n^2 T^2 \) and \( n T \sim P_{\text{SE}}/V_m \), which leads to \( P_{\text{aux}} = R^{0.055} (n/\text{DP})^{1.08} P_0^{0.78} \). Since the alpha and Bremsstrahlung power within \( P_h \) are approximately constant along the \( Q \) lines, the operation map \( P_{\text{aux}}(P_{\text{aux}}, n/\text{DP}) \) (figure 6) is also approximately independent of the reactor size along the constant \( Q \) curve (figure 4). In figure 6 the thermal stability of the operation points is probed by varying the temperature profile by \( 1\% \) and looking at the sign of \( P_h - W/\tau_E \), which is, by construction, equal to zero at each operation point for the equilibrium temperature and density profiles. If a positive (negative) increment in the temperature profile leads to a faster increase (decrease) of the heating power, \( P_h \), compared to the transported power, \( W/\tau_E \), then the operation point is labelled thermally unstable. It is important to note that the shape of the temperature perturbation can affect the sign of the resulting \( P_h - W/\tau_E \). Points below the red curve in figure 6 are at least unstable to a temperature perturbation like the one referred above. While operating a stellarator reactor in thermally unstable conditions might be possible, it would presumably require an active control of the burning point for a stable power output. It should also be noted that the way thermal stability is calculated assumes that the energy confinement scales like equation (3) also ‘locally’, but deviations from it (including confinement transitions) are observed in present individual devices.

4.4. Discussion on the neutron wall loads and breeding technology

As discussed earlier around table 2, smaller reactor devices inevitably suffer more intense neutron bombarding per unit area on the first wall and breeding blanket. To qualify the neutral wall load (NWL) numbers shown in that table we consider that, as a rule of thumb, a maximum NWL of 1.97 MW m\(^{-2}\) (corresponding to a 13.5 m device) would translate into a damage rate of 19.7 dpa/fpy (displacements per atom in a full-power year) in steel and around \( 4 \times 10^{18} \) n/m\(^2\)fpy neutron flux at the first wall. This could be still acceptable under the damage of the first wall, shielding of the other structures (vacuum vessel and coils) and heat recovery points of view, as preliminary assessed in [68], although some improvement on shielding and minor modification on maintenance scheme would be necessary. Nevertheless, such aspects would be not easily manageable with higher NWL. For example, 4.43 MW m\(^{-2}\) \((R = 9 \text{ m in table 2})\) would lead to damage around 40 dpa/fpy and neutron fluxes of almost \( 8 \times 10^{18} \) n/m\(^2\)fpy at the first wall. This would complicate the breeding blanket replacement, since the current Eurofer first wall material suggested to be used in stellarators power plants, as extrapolated from DEMO, is qualified up to 20 dpa for the first DEMO phase (1.57 fpy). Besides, the requirements to the shield components that would need to be very exigent in order to get viable values at the different coil structures On the other hand, the NWL of the largest devices considered in table 2, 0.74 MW m\(^{-2}\) would probably...
lead to low power deposited inside the coolants and accordingly low thermal efficiency. The fourth case with an NWL of 1.11 MW m$^{-2}$, which indeed corresponds to the HSR4/18 design point, could be the best engineering solution to explore since it seems to offer the best compromise between damage/shielding performances, maintenance schemes and heat recovery/thermal efficiency.

Apart from the considerations on the NWL, a substantial increase of the magnetic field can influence other technological aspects related with the in-vessel components. Firstly, important electromagnetic forces can be developed on the structures which form the breeding blanket [69]. It has been demonstrated that the mechanical behaviour of the blanket segments can be compromised. The breeding material can also be affected by the magnetic field. Usually, Li compounds are required to breed the tritium in order to maintain the reactor self-sufficiency. One of the most extended breeding materials is the PbLi eutectic alloy, which indeed is an excellent electric conductor. This means that, when moving inside the blankets, magnetohydrodynamic effects can appear [70]. The main consequence is the increase on the pressure drop of the liquid metal, which has a critical impact on the plant electric efficiency. Moreover, some MHD effects can cause important reverse flows that can compromise the route of the effluent impacting the tritium permeation through the structures or the formation of He clusters. Another important point is corrosion of structural materials due to the interaction with the liquid metal. It has been demonstrated that the presence of an intense magnetic field enhances the corrosion phenomena on EUROFER samples [71].

The conclusion of this discussion is that current breeding blanket technology and maintenance schemes can severely limit the viability of small-size, high-field reactors. The full exploitation of HTS necessitates parallel technological developments on these fronts, that allow one to cope with the increased neutron fluxes. Continuous molten-salt breeders and demountable coil joints, for easier and faster maintenance, have been proposed in the tokamak high-field path to commercial fusion [34]. Some of these ideas have also been considered in stellarator reactor studies [26].

### 4.5. Magnetic field over-engineering

So far we have centred our analysis on design points on the $Q = 40$ line in figure 4. In the 0D analysis of section 3 we argued that increasing the field over those design points, while keeping the $P_{\text{fus}} = 3$ GW target, implied to operate at lower densities and higher temperatures. In a 3D device, operating at lower collisionality and longer alpha particle slowing-down times are disadvantages, as they impose higher standards on the neoclassical and fast particle confinement properties of the magnetic configuration. Since the energy confinement time in its usual form (3) has a positive dependence on the line average density, lowering it is, in a sense, a waste of the confinement potential of the magnetic configurations. However, over-engineering the magnetic field strength in the sense just described brings in two positive consequences of large potential impact: it lowers the plasma $\beta$ and increases the tritium burnup fraction. In reduced-size tokamak reactor studies, high field operation is exploited to move away from operational boundaries [34]. The increase of the magnetic field and the decrease in the operation density both act to lower the critical density fraction $\pi_{c}/\pi_{ec}$. Although it is not in the scope of this article to assess the viability of heat exhaust solutions, we note in passing that those same changes could complicate power handling in the scrape-off layer and divertor regions.

In this subsection we take a look at the consequences of over-engineering the magnetic field with the 1D prescribed-profile analysis utilised in this section. We inspect the effect of a 25% increase in $B$ that for the case of an $R = 13.5$ m device (see table 2). The results are shown in figure 8.

The profiles for the starting $R = 13.5$ m reactor design point are shown in the left column. Those in the centre and right column correspond to the 25% over-engineered field case. In the centre profiles, the ratio of deuterium to tritium densities, $n_{D}/n_{T}$, is kept at 1. In this situation, $n_{D}/n_{T} = 1$, the increase in the fusion power due to the increase of $B$ can only be moderated by reducing the plasma density. Temperature needs to change in roughly inverse proportion to keep up with the 3 GW fusion power. The higher temperatures broaden the
alpha particle generation and heating profile and increase the burnup fraction. The thermal plasma $\beta$ decreases as $B^{-2}$. However, the alpha particle $\beta_\alpha$ slightly increases and broadens as a result of the longer slowing-down time and the broader profile of alpha particle production. These two effects would demand better fast particle confinement properties (in particular for the longer time-scale stochastic diffusion losses) in a more extended radial range for this over-engineered $B$-field case with $n_D/n_T = 1$. It should also be noted that for HELIAS-type configurations, a sufficient plasma beta is required to reduce fast ion prompt losses, but the necessary minimum beta is subject to some adjustment at the stage of the optimization of the magnetic configuration. The last row of figure 8 shows the comparison of the normalised thermal diffusivity $\chi$ (see appendix A.2) with the proxy for the neoclassical $1/\nu$ diffusivity. The effective ripple is set to $\epsilon_{\text{eff}} = \epsilon_{\text{eff}} - 0.5\%$. These plots illustrate that further reductions of $\epsilon_{\text{eff}}$ could be necessary to make the core profiles compatible with the neoclassical heat transport levels in the over-engineered case with $n_D/n_T = 1$.

Instead of lowering the plasma density, in the right column of figure 8 the ratio of deuterium to tritium has been increased to moderate the fusion power. This is shown to result in a moderation of the negative consequences referred to above, that stemmed from the necessary decrease of density and increase of temperature. While the reduction of the plasma $\beta$ is not as pronounced, the $\beta_\alpha$ is also reduced and the collisionality increases only slightly (bottom plot). The reduction of the tritium concentration increases the burnup fraction three-fold with respect to the normal field case. We conclude that magnetic field over-engineering could be taken advantage of by reducing the tritium concentration, thereby allowing to reduce the plasma and alpha particle normalised pressures and increase the burnup fraction, without strongly reducing plasma collisionality.

4.6. Design points of stellarator reactors with constant neutron wall loading

As mentioned before, the NWL is an important technological parameter. The breeding blanket and heat-to-electricity transformation technology determine an optimal range for the NWL. Therefore, one could be interested to search for reactor design points of constant NWL rather than constant fusion power. Reducing the size of the reactor design point while keeping a constant NWL still gives favourable reduction of the ratio of GW per m$^3$ of reactor material, but the larger electro-mechanical stresses of smaller, higher-field devices could make the support structure bigger and more expensive than that implied by a simple $R^3$ scaling.
Referring to the operation map in figure 6, valid for all reactor design points in table 2, one sees that reducing the auxiliary heating power and adapting the density, one can go down the $Q = 40$ line to the fusion power compatible with a given reduced NWL. However, in doing so, one would require the reactor to operate at larger ratios of density over the critical density $\rho_c/\rho_{ec}$, since $\rho_{ec} \sim P_f/\rho_i$. In consequence, if $\rho_c/\rho_{ec}$ is to be kept constant while reducing the fusion power at fixed device size, the magnetic field needs to be increased. This is a consequence of the strong power dependence of the critical density characteristic of stellarators [72]. The resulting $(B, R)$ design points and some of their physics characteristics are shown in figure 9 and table 3 respectively. Besides the critical density fraction, the fusion gain is kept constant at $Q = 40$ in those points. As shown in table 3, smaller, higher-field devices with constant NWL (i.e. $P_{las} \sim R^{-2}$) display a rapidly decreasing plasma beta, whereas the central collisionality decreases more slowly. The plasma temperature varies only slightly in the explored range.

It should be emphasised that the present understanding of the Sudo-type density limit in stellarators does not allow to determine an absolute limit as is the case for tokamaks’ Greenwald limit [72]. In this study we have used the published working point of the HSR4/18 stellarator to define the reference $\rho_c/\rho_{ec}$, but the possibility of operating at substantially higher densities cannot be ruled out. If this were to be the case, the magnetic field of the design points could be lowered with respect to those listed in table 3. According to the present understanding (see e.g. [73] and references therein), the density limit in stellarators is connected to the properties and concentration of the edge radiative impurity. Gaining a more predictive capability for the determination of the critical density in reactor conditions is therefore linked with the validation of models of divertor and SOL/edge impurity transport and with the constraints on the concentration of impurities imposed by power exhaust.

### 5. Summary and conclusions

In this article we have shown that the design points of stellarators of different scale and field, but otherwise similar fusion power and gain, share many similarities under the assumption that the plasma density design point can be scaled as $n \sim B^2$. Archetypical profiles for temperature, plasma $\beta$, gradient scale lengths or thermal diffusivity in gyro-Bohm units, among others, characterise a family of varying-field stellarator reactors with maximally-reduced size, $R \sim B^{-4/3}$ (figure 5). The suitability requirements on the magnetic configuration (e.g. good flux surfaces and MHD stability at high-$\beta$, reduced neoclassical transport in the low-collisionality regime or sufficient confinement of alpha particles) are therefore largely independent of the device size/field strength across the family of reactor design points. In this sense, the qualification of optimization criteria in devices like the W7-X is still relevant in a future scenario in which electromagnets based on HTS can be applied in stellarator reactors. Furthermore, we have shown that the operation landscape of fusion power and gain as function of the auxiliary power and relative density is also shared among the family of reactors (figure 6). In consequence, if the technological development of high-field-capable electromagnets progresses to make ten-Tesla-class stellarators accessible, a demonstration device of much reduced size ($<10$ m) is conceivable, that would qualify larger-scale devices in several meaningful ways.

High field and reduced size make conditions particularly harsh for the breeding structure, which calls for alternative approaches for breeder and wall maintenance [34]. An additional complication for devising reduced size stellarator reactor stems from the need to increase the distance from the magnets to the plasma edge (relative to the plasma radius) to allow for the irreducible space of the breeding blanket and neutron shield. This results in an increase in the complexity of the coils that would be required to generate an optimised magnetic configuration [22]. It is in this respect that higher-field operation would call for the search of optimised configurations with a larger relative distance between the current and control surfaces. The ability to relax some of the optimization targets, based on possibly over-simplified physics models, would make this line appear more promising.

If the tritium breeding technology or limitations in remote maintenance were to set the minimum size of a stellarator reactor, one could wish to increase the field strength while maintaining a certain device size. Magnetic field over-engineering is accompanied by reduced thermal $\beta$ and higher tritium burnup fraction, but conditions on thermal and alpha particle confinement become more stringent. However, these drawbacks can be largely mitigated with the adjustment of the D/T density ratio (figure 8).

As an alternative to the constant-$P_{las}$ scaling of reactor design points adopted throughout this article, the scaling at constant NWL and gain has also been inspected in this article. The stellarator-specific critical density dependence on heating power results in a stronger field scaling ($R \sim B^{-1.15}$) when the reactor linear size is reduced at constant NWL and critical density fraction (figure 9). The resulting reactor design points feature lower plasma $\beta$ and slightly lower collisionality compared to same size, higher $P_{las}$ design points.

### Table 3. Reactor design points of constant NWL = 1.11 MW m$^{-2}$, $\rho_c/\rho_{ec}$ and $Q = 40$ shown in figure 9.

| $R$ (m) | $B$ (T) | $\rho_c(10^{20}$ m$^{-3}$ | $(T)$ (keV) | $P_{las}$ (GW) | $\beta$ (%) | $\beta(0)/\beta_c(0)$ | $\nu^2(0)(10^{-3})$ | $(\chi/\chi_{ph})_{0.5}$ |
|--------|--------|------------------------|-------------|---------------|-------------|----------------------|---------------------|---------------------|
| 7.00   | 14.58  | 3.89                   | 5.54        | 0.45          | 0.79        | 3.68                 | 0.77                | 0.05                |
| 9.00   | 10.86  | 3.58                   | 5.54        | 0.75          | 1.27        | 3.91                 | 0.96                | 0.06                |
| 13.50  | 6.89   | 2.98                   | 5.25        | 1.69          | 2.58        | 4.04                 | 1.23                | 0.07                |
| 18.00  | 5.00   | 2.60                   | 5.21        | 3.00          | 4.24        | 4.10                 | 1.44                | 0.07                |

Nucl. Fusion 62 (2022) 036024
The conclusions of this work depend on a number of assumptions that have been presented but cannot be justified on solid physics grounds. The close invariance of plasma profiles and operation map (figures 5 and 6) depends on the use of prescribed shapes for the density and temperature profiles, that are assumed to be the same for all reactors. Other assumptions such as the fuel-mix composition and the alpha heating efficiency add quantitative uncertainty to the design points that have been considered. Our analysis critically relies on the ISS04 scaling of the energy confinement time. The assessment of any limitation to this scaling in high-beta, low-collisionality regimes should be regarded a stellarator research priority. Finally, it should not be left unnoticed that the viability of power exhaust has not been considered in any respect in our study. The development of stellarator scrape-off layer and divertor physics models that allow to incorporate exhaust conditions to reactor studies also appears to be a necessary step in the development of a credible stellarator reactor concept. Significant departures with respect to the tokamak results [6, 7] could arise from the absence of a well defined threshold power across the separatrix and from differences in the scaling of the scrape-off layer width [13, 74]. This development bears also the importance of allowing to quantify the Sudo-type density limit that is compatible with exhaust conditions in stellarator reactor studies.

Acknowledgments

We are thankful to C.D. Beidler, S. Bozhenkov, H. Laqua, T. Sunn Pedersen, F. Warmer, R. Wolf and H. Zohm for reading the manuscript and providing useful comments. We thank F. Parra for discussions on the content of section 4.1.2, P. Helander for suggesting the study of increased D/T fraction of section 4.5 and J. de la Riva for his help implementing changes in the code. The first author acknowledges the sustained interaction with C.D. Beidler, A. Dinklage and F. Warmer on stellarator physics and reactor studies. This work is partially funded by the research Grant PGC2018-095307-B-100 from the Spanish Ministerio de Ciencia, Innovación y Universidades.

Appendix A. Definitions and conventions

A.1. Model plasma profiles

Missing a complete model for particle and energy sources and transport, we fix the shape of the density and temperature profiles used for the 1D analysis. The density profiles is chosen to be represented by

\[ \hat{n}_{\text{model}}(\rho) = \pi \frac{1}{2} - \text{atan} \left( k_1 (\rho^2 - \rho_0^2) \right) - k_2 (\rho^2 - 1), \]  
\[ \text{with } \rho_0^2 = 0.95 \text{ (location of the edge density gradient)}, \]  \[ k_1 = 30 \text{ (steepness of the gradient)} \]  \[ k_2 = 0.5 \text{ (core flatness)}. \]

Rather than fixing the form of the temperature profile, we choose to make it dependent on the profile of the net heating power density. We first define the profile function,

\[ F(r) = -\int_{0}^{r} \frac{1}{r} \sum_{i} n_i(r') \left( \int_{0}^{r'} r'^2 S(r') dr'' \right) dr', \]  
with \( S(r) = k_o S_\alpha + S_{\text{aux}} - S_B \). The model temperature profile is expressed in terms of \( F(r) \),

\[ T_{\text{model}}(r) = 1 + k - F(r). \]  
where \( k = 0.01 \) relates to the core to edge temperature ratio, \( T(\alpha)/T(0) = k/(1 + k) \). The profile given by (A.3) can be seen to result in a radially-constant heat diffusivity.

A.2. Heat diffusivity

The total thermal diffusivity is defined as the weighted average of the individual species diffusivities \( \chi = \sum_i \chi_i n_i / \sum_i n_i \). The total heat flux \( \Gamma_Q \) is written as \( \Gamma_Q = -\frac{1}{2} \sum_i n_i \frac{dF}{dr} \). The steady-state heat balance equation then relates this flux to the heat density \( S = k_0 S_\alpha + S_{\text{aux}} - S_B \),

\[ \int \frac{1}{r} \frac{d}{dr} r \Gamma_Q = S. \]  
In some of the plots this is normalised by the local gyro-Bohm diffusivity, defined by \( \chi_{\text{gb}} = \alpha v_T (\rho_s)^2 \), where the subindex T refers here to tritium as a plasma species.

A.3. Normalised gyroradius and plasma collisionality

The normalised gyroradius of a plasma species \( s \) is defined in this article as

\[ \rho_s^* = \frac{\nu_s}{\Omega_b}, \]  
where \( \nu_s \) is the thermal velocity \( \nu_s = \sqrt{2 T_s / m_s} \) and \( \Omega_b \) is the gyrofrequency \( \Omega_b = Z_e e B / m_s \).

As to the collision frequency of each species, \( \nu_c \), we adopt the definition given in [63, page 3]. Collisionality is then given by \( \nu_c^* = \nu_c R / v_s \). For clarity, we only discuss triton collisionality in the main text, which accounts for self-collisions and collisions with deuterons. For equal temperatures, triton collisionality is slightly smaller than deuteron collisionality and about a factor 3 smaller than electron collisionality.

A.4. α-particle pressure

To estimate the equilibrium alpha particle pressure and beta we use the classical slowing-down distribution function (see e.g. [75])

\[ f_{\alpha}(v) = \frac{s_\alpha \tau_{\alpha}}{4\pi (v^3 + v_c^3)} H(v_s - v), \]  
where \( s_\alpha = n_\alpha \rho_T \langle \sigma v \rangle / m_\alpha = \frac{1}{2} \langle \sigma v \rangle / m_\alpha \langle \rho_T \rangle \) is the production rate of alphas per unit volume with a birth velocity \( v_s \), and \( H \) is the Heaviside function.
function. The slowing-down time $\tau_s$ is defined as

$$
\tau_s^{-1} = \frac{4}{3\sqrt{n_0}} \frac{Z^2 e^2}{3} n_0 \ln \Lambda,
$$
(A.8)

whereas the critical velocity $v_c$ is given by

$$
v_c = \left( \frac{mZ^2 e^2}{3} \frac{\sqrt{n_0}}{m_i} \right)^{1/3} v_{te},
$$
(A.9)

In this equation we consider the deuterium and tritium a single species with 2.5 amu.

The alpha particle pressure is then given by

$$p_a = \int d^3v \frac{m_a v^2}{3} f_{\alpha}(v) = \frac{m_a v^2}{3} s_\alpha s_e I \left( \frac{v_0}{v_c} \right),
$$
(A.10)

and

$$I(x) = \int_0^x \frac{dt}{t^4 + 1}.
$$
(A.11)

The alpha particle beta, $\beta_\alpha = p_a/(B^2/2\mu_0)$ is a weak function of the reactor size along the $Q = 40$ line in figure 4. In fact, assuming a constant temperature (equation (12)), the critical velocity (A.9) is constant and (A.10) scales like $p_a \sim s_\alpha \tau_s \sim n$. With the scaling (11), this results in $\beta_\alpha \sim n/B^2 = $ constant.

A.5. Tritium burnup fraction

The so-called tritium burnup fraction is a global magnitude defined as the ratio between the rate of fusion reactions within the confined volume and the fuelling rate of fresh tritium (see e.g. [76]), namely

$$f_{\text{burnup}} = \eta_{\text{eff}} \left( 1 + \frac{1}{2 \tau_s} \right),
$$
(A.12)

where $\eta_{\text{eff}}$ is the fuelling efficiency (which depends, in particular, on the fuelling method) and $N_T = \int_0^x dV n_T$ and $K_a = \int_0^x dV n_I \sigma(v) n_T$. The effective tritium confinement time, $\tau_T$, is defined in terms of the tritium inventory and losses in a power plant. The figure of merit shown in the figures corresponds to (A.12) with $\eta_{\text{eff}} = 1$, $R = 0$, $\tau_T = \tau_E$ and $N_T = N_D = 0.45N_e$. In figure 3 we plot a related quantity labelled tritium burnup profile, that is defined as $k_T = n_T \sigma(v) n_T \tau_T$.

With this definition and assumptions, the tritium burnup fraction is almost constant along the reactor line in figure 4 ($P_{\text{fus}} = 3$ GW, $Q = 40$). Since the fusion power is kept constant and is proportional to the alpha power, the $K_a = P_a/E_a$ is constant. The ratio $\tau_E/(N_D + N_T) \sim \tau_E/(nV_a)$ then scales like plasma temperature in equation (12).

Appendix B. Defining the operation point for density

The choice of density operation point for a reactor bears considerable importance. The fusion power scales approximately as $P_{\text{fus}} \sim n^2 T^2 V_a$, so that, for a fixed fusion power, the required temperature scales inversely with density. While the required $\beta$ is roughly independent of the density operation point, the characteristic collisionality ($\nu_s \sim nT^2$) increases quickly with density. The energy confinement time has a positive dependence on density, whereas the alpha particle slowing-down time decreases with it. Furthermore, the divertor operation might require high density operation together with extrinsic radiators in the divertor and burnup region for power exhaust. One concludes that the density design point should be made as high as allowed by operation limits and heating/fuelling access.

In this study, we have chosen to scale density as $B^2$, which is consistent with the increase of the ECRH cut-off density (see the lower plot in figure B1). We note that the chosen DP for density is not compatible with having an O1 ECRH heating scheme. Second harmonic fast-X mode or first harmonic slow-X mode (with high-field-side launching) would be required. The gyrotron frequencies that are needed for high field operation have not been achieved to date, so that a technological development would be necessary on this front for operating high field stellarators. In this article we adopt the view that the heating method, whether ECRH or others, should not be the primary density-limiting factor and that such technological developments can be carried out.

The top plot of figure B1 shows that the chosen density scaling is compatible with keeping a constant ratio of density DP to the radiative density limit. The original Sudo density limit [77] for the LHD stellarator is $n_c \sim \sqrt{P_B}/V_a$ similar to the one found in W7-AS [78]. More recently, the W7-X density limit has been shown to be in agreement with a similar scaling [73], namely

$$n_c \propto P_h^{0.57} B^{0.34} R^{-1.25},
$$
(1.1)

where a constant aspect ratio has been assumed ($a = R/A$). This limit is understood to be due to the critical temperature behaviour of edge impurity radiation [79], which relates the density limit with the quality of the confinement and the specific characteristics of the radiator that are accounted for in the proportionality factor. The question whether such a limit is to be expected in a reactor environment, where the current edge radiators, carbon and oxygen, will not be present is in order. Tungsten as the main wall element displays a cooling factor with only mild temperature dependence. However, it is to be expected that tungsten levels are kept very low and extrinsic edge radiators are used to control power exhaust in the edge and divertor regions. These would also exhibit higher cooling rates below a critical temperature and, in this sense, a density limit scaling of the Sudo type could be relevant for reactor conditions.
Taking the scalings (11), (13) and (B.1) it easy to show that

$$\frac{n}{n_c} \sim P_i^{0.57} B^{-0.01}.$$  \hfill (B.2)

So, for constant fusion power and gain and a density scaling as $n \sim B^2$, the ratio $n/n_c$ remains constant as shown in figure B1. This density scaling is therefore consistent with the presently known stellarator density limit.

**References**

[1] Abdou M., Riva M., Ying A., Day C., Loarte A., Baylor L.R., Hamrickhouse P., Fuerst T.F. and Cho S. 2020 Physics and technology considerations for the deuterium–tritium fuel cycle and conditions for tritium fuel self-sufficiency Nucl. Fusion 61 013001

[2] Dinklage A. et al 2007 Assessment of global stellarator confinement: status of the international stellarator confinement database Fusion Sci. Technol. 51 1–7

[3] Erckmann V. et al 1993 H mode of the W 7-AS stellarator Phys. Rev. Lett. 70 2086–9

[4] Hirsch M. et al 2008 Major results from the stellarator Wendelstein 7-AS Plasma Phys. Control. Fusion 50 055001

[5] Estrada T. et al 2010 L–H transition experiments in TJ-II Contrib. Plasma Phys. 50 501–6

[6] Reinke M.L. 2017 Heat flux mitigation by impurity seeding in high-field tokamaks Nucl. Fusion 57 034004

[7] Siccinio M., Fable E., Angioni C., Saarelma S., Scarabosio A. and Zohm H. 2017 Impact of an integrated core/SOL description on the R and BT optimization of tokamak fusion reactors Nucl. Fusion 58 016032

[8] Ohyabu N. et al 1994 The large helical device (LHD) helical divertor Nucl. Fusion 34 387–99

[9] Boozer A.H. 2015 Stellarator design J. Plasma Phys. 81 515810606

[10] Bader A., Boozer A.H., Hegna C.C., Lazerson S.A. and Schmitt J.C. 2017 HSX as an example of a resilient non-resonant divertor Phys. Plasmas 24 032506

[11] Sardei F. et al 1997 Island divertor studies on W7-AS J. Nucl. Mater. 241–243 135–48

[12] Renner H., Sharma D., Kühlinger J., Boscardy J., Grote H. and Schneider R. 2004 Physical aspects and design of the Wendelstein 7-X divertor Fusion Sci. Technol. 46 318–26

[13] Niemann H. et al 2020 Large wetted areas of divertor power loads at Wendelstein 7-X Nucl. Fusion 60 084003

[14] Winters V.R. 2019 Carbon sourcing and transport in the island divertor of Wendelstein 7-X PhD Thesis University of Wisconsin–Madison (https://search.library.wisc.edu/catalog/991289086902121)

[15] Zhang D. et al 2019 First observation of a stable highly dissipative divertor plasma regime on the Wendelstein 7-X stellarator Phys. Rev. Lett. 123 025002

[16] Schmitz O. et al 2020 Stable heat and particle flux detachment with efficient particle exhaust in the island divertor of Wendelstein 7-X Nucl. Fusion 61 016026

[17] Mayer M. et al 2020 Material erosion and deposition on the divertor of W7-X Phys. Scr. T171 014035

[18] Killer C., Shanahan B., Gruhke O., Endler M., Hammond K. and Rudischhauser L. 2020 Plasma filaments in the scrape-off layer of Wendelstein 7-X Plasma Phys. Control. Fusion 62 085003

[19] Feng Y. et al 2021 Understanding detachment of the W7-X island divertor Nucl. Fusion 61 086012

[20] Brezinsek S. et al 2021 Plasma-surface interaction in the stellarator W7-X: conclusions drawn from operation with graphite plasma-facing components Nucl. Fusion 62 016006

[21] Warmer F., Bykov V., Drevlak M., Häußler A., Fischer U., Stange T., Beidler C.D. and Wolf R.C. 2017 From W7-X to a...
HELIAS fusion power plant: on engineering considerations for next-step stellarator devices Fusion Eng. Des. 123 47–53

[22] Landreman M. and Booher A.H. 2016 Efficient magnetic fields for supporting toroidal plasmas Phys. Plasmas 23 032506

[23] Yamaguchi H. 2019 A quasi-isodynamic magnetic field generated by helical coils Nucl. Fusion 59 104002

[24] Najmabadi F. et al. 2008 The ARIES-CS compact stellarator fusion power plant Fusion Sci. Technol. 54 655−72

[25] Goto T., Miyazawa J., Tanaka T., Yanagi N. and Sago R. 2011 A robust design window for the heliotron DEMO reactors Plasma Fusion Res. 6 2405083

[26] Sagara A., Miyazawa J., Tamura H., Tanaka T., Goto T., Yanagi N., Sakamoto R., Masuzaki S. and Ohtani H. 2017 Two conceptual designs of helical fusion reactor FFHR-d1A based on ITER technologies and challenging ideas Nucl. Fusion 57 086046

[27] Grieger G. et al. 1992 Modular stellarator reactors and plans for Wendelstein 7-X Fusion Technol. 21 1767−78

[28] Beidler C.D. et al. 2001 The HELIAS reactor HSR/418 Nucl. Fusion 41 1759−66

[29] Warmer F., Beidler C.D., Dinklage A., Turkin Y. and Wolf R. 2015 Limits of confinement enhancement for stellarators Fusion Sci. Technol. 68 727−40

[30] Beidler C.D., Hamreyer E., Herrnegger F., Kisslinger J., Igitkhanov Y. and Wobig H. 2001 Stellarator fusion reactor-an overview (Toki, Gifu (Japan) 11-14 December 2001) Toki Conf. ITC12 vol 30p 2014

[31] Lion J., Warmer F., Wang H., Beidler C.D., Muldrew S.I. and Wolf R.C. 2021 A general stellarator version of the systems code PROCESS Nucl. Fusion 61 126021

[32] Fisk Z., Thompson J.D., Zirngiebel E., Smith J.L. and Cheong S.-W. 1987 Superconductivity of rare earth-barium–copper oxides Solid State Commun. 62 743−4

[33] Sorbom B.N. et al. 2015 ARC: a compact, high-field, fusion research nuclear facility concept and demonstration power plant with demountable magnets Fusion Eng. Des. 100 378−405

[34] Whyte D.G., Minervini J., LaBombard B., Marmar E., Xanthopoulos P., Heidbrink W.W. 2008 Basic physics of Alfvén instabilities driven by energetic particles in toroidally confined plasmas Phys. Plasmas 15 055501

[35] Di Siena A., Bañez J.A. 2008 Optimization of energy confinement in optimized stellarators Nucl. Fusion 58 063022

[36] Bozhenkov S.A. et al. 2020 High-performance plasmas after pellet injections in Wendelstein 7-X Nucl. Fusion 60 066011

[37] Bakam M., Pogn F.I. and Landreman M. 2019 stella: an operator-split, implicit-explicit 4f–gyrokinetic code for general magnetic field configurations J. Comput. Phys. 391 365−80

[38] Mauer M. et al. 2020 GENE-3D: a global gyrokinetic turbulence code for stellarators J. Comput. Phys. 420 109604

[39] Xanthopoulos P. et al. 2020 Turbulence mechanisms of enhanced performance stellarator plasmas Phys. Rev. Lett. 125 075001

[40] Heidbrink W.W. 2008 Basic physics of Alfvén instabilities driven by energetic particles in toroidally confined plasmas Phys. Plasmas 15 055501

[41] Di Siena A., Bañez J.A. et al. 2020 Turbulence suppression by energetic particle effects in modern optimized stellarators Phys. Rev. Lett. 125 105002

[42] Nishikawa M. 2010 Study on tritium balance in a D−T fusion reactor Fusion Sci. Technol. 57 120−8

[43] Nishikawa M. 2011 Tritium balance in a D−T fusion reactor Fusion Sci. Technol. 59 150−62

[44] Tamor S. 1988 Synchrotron radiation loss from hot plasma Nucl. Instrum. Methods Phys. Res. A 271 37−40

[45] Helder P., Parra F.I. and Newton S.L. 2017 Stellarator bootstrap current and plasma flow velocity at low collisionality J. Plasma Phys. 83 095803026

[46] Boozer A.H. and Gardner H.J. 1990 The bootstrap current in stellarators Phys. Fluids B 2 2408−21

[47] Sewald B., Kaslov S.V., Kernbichler W., Kaluzhnyj V.N., Nemov V.V., Tribaldos V. and Jiménez J.A. 2008 Optimization of energy confinement in the 1/4 regime for stellarators J. Comput. Phys. 227 6165−83

[48] Beidler C.D. et al. 2011 Benchmarking of the mono-energetic transport coefficients results from the International Collaboration on Neoclassical Transport In Stellarators (ICNTS) Nucl. Fusion 51 076001

[49] Hemmerberg S.A., Drevlak M., Nührenberg C., Beidler C.D., Turkin Y., Loizu J. and Helder P. 2019 Properties of a new quasi-axisymmetric configuration Nucl. Fusion 59 026014

[50] Nakata M., Nunami M., Sugama H. and Watanabe T.-H. 2017 Ionotope effects on trapped-electron-mode driven turbulence and zonal flows in helical and tokamak plasmas Phys. Rev. Lett. 118 165002

[51] Beidler C.D., Kolosnichenko Y.I., Marchenko V.S., Sidorenko I.N. and Wobig H. 2001 Stochastic diffusion of energetic ions in optimized stellarators Phys. Plasmas 8 2731−8
[67] Beidler C.D., Feng Y., Geiger J., Küchl F., Maßberg H., Marushchenko N.B., Nührenberg C., Smith H.M. and Turkin Y. 2018 (Expected difficulties with) density-profile control in W7-X high-performance plasmas Plasma Phys. Control. Fusion 60 105008

[68] Palermo I., Warner F. and Hußler A. 2021 Nuclear design and assessments of helical-axis advanced stellarator with dual-coolant lithium-lead breeding blanket; adaptation from DEMO tokamak reactor Nucl. Fusion 61 076019

[69] Maione I.A., Roccella M., Marin A., Bertolini C. and Lucca F. 2019 A complete EM analysis of DEMO WCLL breeding blanket segments during VDE-up Fusion Eng. Des. 146 198–202

[70] Urgorri F.R., Smolentsev S., Fernández-Berceruelo I., Rapisarda D., Palermo I. and Ibarra A. 2018 Magneto-hydrodynamic and thermal analysis of PbLi flows in poloidal channels with flow channel insert for the EU-DCLL blanket Nucl. Fusion 58 106001

[71] Gómez C., Hernández T., Muktpavela F., Platacis E. and Shishko A. 2015 Magnetic field effect on the corrosion processes at the Eurofer–Pb–17Li flow interface J. Nucl. Mater. 465 633–9

[72] Greenwald M., Hernández T., Muktpavela F., Platacis E. and Shishko A. 2002 Density limits in toroidal plasmas Plasma Phys. Control. Fusion 44 R27

[73] Fuchert G. et al 2020 Increasing the density in Wendelstein 7-X: benefits and limitations Nucl. Fusion 60 036020

[74] Eich T. et al 2013 Scaling of the tokamak near the scrape-off layer H-mode power width and implications for ITER Nucl. Fusion 53 093031

[75] Helander P. and Sigmar D.J. 2002 Collisional Transport in Magnetized Plasmas (Cambridge Monographs on Plasma Physics vol 4) (Cambridge: Cambridge University Press)

[76] Jackson G.L., Chan V.S. and Stambaugh R.D. 2013 An analytic expression for the tritium burnup fraction in burning-plasma devices Fusion Sci. Technol. 64 8–12

[77] Sudo S., Takeiri Y., Zushi H., Sano F., Itoh K., Kondo K. and Iiyoshi A. 1990 Scalings of energy confinement and density limit in stellarator/heliotron devices Nucl. Fusion 30 11–21

[78] Giannone L. et al 2000 Physics of the density limit in the W7-AS stellarator Plasma Phys. Control. Fusion 42 603–27

[79] Itoh K. and Itoh S.-I. 1988 Detached and attached plasma in stellarators J. Phys. Soc. Japan 57 1269–72