Neutral beam current drive in a tokamak

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Abstract
Neutral beam current drive (NBCD) on the EAST tokamak is studied by using Monte-Carlo simulation. The electron shielding effect to the fast ion current is taken into account by using a fitting formula applicable to general tokamak equilibria and arbitrary collisionality regime. The net currents driven by the beam are off-axis although the fast ion currents are on-axis. This is found to be due to the electron shielding effect being strong near the magnetic axis. We also investigate the dependence of NBCD efficiency on the plasma density. The results indicate that the NBCD efficiency decreases with the increase in plasma density. A simple semi-analytic estimation of the dependence of NBCD efficiency on the density is proposed and is in reasonable agreement with the results directly obtained in the simulations.

1 Introduction
Neutral beam injection (NBI) is widely used in tokamaks for heating plasma, driving plasma rotation, and driving electric current. It is straightforward to compute the steady-state current carried by NBI fast ions using Monte-Carlo test particle simulations. However, to get the net current, one must take into account the electron shielding effect, i.e., the current carried by electrons due to the response of electrons to the presence of fast ions. This generally requires to solve the steady-state Fokker-Planck equation for electrons with additional collision term corresponding to the electron collision with the fast ions. For current drive problem, we are interested in the first (parallel) moment of the electron distribution function. Then, making use of the self-adjoint property of the linearized collision operator, the electron response to arbitrary fast ion sources can be obtained by using the Green function method. The final results of these studies are usually some fitting formulas for the ratio of net current to the pure fast ion current. The present work uses these fitting formulas to include the electron shielding effect. The electron shielding model used in this work is a general model applicable to arbitrary collisionality regime and general tokamak flux surface shapes.

Using the electron shielding model mentioned above and a Monte-Carlo test particle code, we studied the neutral beam current drive (NBCD) on the EAST tokamak. The results indicates the net current is often off-axis even when the pure fast ion current is on the axis. This is found to be due to electron shielding effect being strong near the magnetic axis.

We also study the dependence of the neutral beam current drive (NBCD) efficiency on the plasma density. Define NBCD efficiency by

\[ \eta = \frac{I_{cd}}{P_{nbi}}, \]

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where $I_{cd}$ is the net current driven by the neutral beam (fast ion current minus the electron shielding current), and $P_{nbi}$ is the beam source power after leaving the neutralization vessel, which is fixed at 1 MW in the simulation.

The simulations indicate that the NBCD efficiency decreases with increase in the plasma density. The drive efficiency reaches a large value in a very low density, even though the shine-through loss is significant there. With the increase of the density, the drive efficiency decreases, which is found to be mainly due to the decreasing of the slowing-down time. This is because, when the slowing-down time becomes shorter, fast ions can remain energetic for shorter time and thus less fast ions can contribute to the fast ion current. A simple semi-analytical formula is proposed to take into account of the two competing factors relating to the current drive, namely, the loss fraction and slowing-down time. The dependence of NBCD efficiency on the density given by the formula is in reasonable agreement with the simulations.

We use the guiding-center drift model in simulating the fast ion behavior. The finite Larmor radius (FLR) effects of fast ions can be included in the simulations when necessary. We compare the results obtained when the FLR effects are included with those obtained when they are not. The difference between the two cases is negligible, indicating the FLR effects are not important for the parameter regime studied in the work.

## 2 Plasma equilibrium configuration and profiles

The simulations are performed for the EAST tokamak, which is a superconducting tokamak with a major radius $R_0 = 1.85 m$, minor radius $a \approx 0.45m$, typical on-axis magnetic field strength $B_0 \approx 2.2T$ and plasma current $I_p \approx 0.5 MA$\cite{13, 14}. The magnetic configuration and plasma profiles used in this work (Fig. 1) were reconstructed by the EFIT code\cite{16} from EAST tokamak discharge #102392@3490ms with constrains from experiment diagnostics.

![Figure 1: Left panel: profiles of electron number density $n_e$, electron temperature $T_e$, ion temperature $T_i$, and safety factor $q$. The radial coordinate $\rho$ is the square root of the normalized poloidal magnetic flux: $\rho_p = \sqrt{(\Psi - \Psi_0)/(|\Psi_b - \Psi_0|)}$, where $\Psi \equiv A_p R$ is the poloidal flux function, $\Psi_0$ and $\Psi_b$ are the values of $\Psi$ at the magnetic axis and last-closed-flux-surface, respectively. Right panel: magnetic configuration. This is a low single null configuration with $B_{axis} = 2.2T$, $I_p = 404 kA$, $q_{axis} = 0.95$, and $q_{95} = 5.05$. Directions of the current and toroidal magnetic field are indicated in the figure.](image)
In this work, for simplicity, we consider a pure Deuterium plasma, and thus $Z_i = Z_{\text{eff}} = 1$ and $n_i = n_e$, where $Z_i$ and $n_i$ are charge number and number density of background ions, respectively, $Z_{\text{eff}}$ is the effective charge number of background ions.

### 3 Simulations

Monte-Carlo simulations in this work are performed by using TGCO code\[^{13, 17}\], which models neutral beam ionization and collision transport of the resulting fast ions. We consider a co-current Deuterium NBI with tangential radius $R_{\text{tan}} = 1.141\, m$, full energy $E_{\text{full}} = 55\, \text{keV}$, and the particle number ratio between full, half, and $1/3$ energy being 80% : 14% : 6%. More details about the fast ion source are provided in Appendix A.

The relaxation process after the beam is turned on is shown in Fig. 2, where the time evolution of the fast ion current and the resulting net current are plotted. The results show that the current saturates in about 50ms, which is about the typical fast ion slowing-down time for these parameters.

The process of the electron shielding the fast ion current in the relaxation is not directly modeled by the Mont-Carlo code. The shielding effect is taken into account only by multiplying the fast ion current by a shielding factor $F$ (the details are given in 3). We are interested in the steady-state current, which should be independent of the behavior of electron shielding in the relaxation process. Results in Fig. 2 show that the steady-state net current is about 1/3 of the fast ion current and the resulting current drive efficiency is about 10 kA / MW.

![Figure 2: The relaxation process of fast ion current and net current after the beam is turned on. Beam power is fixed at 1 MW.](image)

The steady-state fast ion current is obtained from integrating the fast ion distribution in phase space. Figure 3 plots the steady-state fast ion distribution in $(E, v_\parallel/v)$, where $E$ is the kinetic energy and $v_\parallel$ is the parallel (to the magnetic field) velocity. There are three jumps in Fig 3(a) and (b), which correspond to the NBI source at full, half, and 1/3 of 55 keV. Note that collisional energy diffusion makes some particles exceed the full energy, as is shown in Fig 3(b). The average particle kinetic energy obtained in the simulation is 26 keV.

Figure 4 plots the radial profiles of various quantities related to the driven current, namely, the pure fast ion current density $J_f$, the electron shielding factor $F$, the net current density
Figure 3: Steady-state distribution of NBI fast ions in \((E,v_∥/v)\) (panel (a)), in \(E\) (panel (b)), and in \(v_∥/v\) (panel (c)). Here \(f_E\) is defined by \(dN = f_E dE\), where \(dN\) is the number of particles within the energy interval \(dE\). And \(f_{v_∥/v}\) is defined in a similar way, i.e., \(dN = f_{v_∥/v} d(v_∥/v)\), where \(dN\) is the number of particles within the interval \(d(v_∥/v)\). The critical energy for this case \(E_{\text{crit}} = 42.4\) keV.

\[ J_{\text{net}} = J_f F, \text{ the effective trapped electron fraction } f_t, \text{ the electron collision frequency } \nu_e, \text{ the thermal electron bounce frequency } \omega_b, \text{ and the normalized electron collision frequency } \nu_{e\star}. \] (The details of these quantities are given in Appendix C.) The formulas for the shielding effect used here are valid for general tokamak equilibria and arbitrary collisionality regime[11, 12], where the equilibrium shaping effects are included via the effective trapped fraction \(f_t\).

Figure 4 shows that the pure fast ion current peak toward the magnetic axis, i.e., it is on-axis. However, the net current reach its peak at \(\rho_p \approx 0.3\), i.e, it is off-axis. This is due to the electron shielding effect, which, as is shown by the factor \(F\) approaching zero near the magnetic axis, is strong near the magnetic axis. The reason for \(F\) to reach its peak value at an off-axis location is as follows. The formula for the \(F\) factor (as is detailed in Appendix C) depends on two factors, namely, effective trapped fraction \(f_t\) and the normalized electron collision frequency \(\nu_{e\star}\). From the core toward the edge, the effective trapped fraction is increasing and thus less electrons can contribute in canceling the fast ion current. Therefore we see \(F\) increases with the radial.
Figure 4: Radial profiles of effective trapped electron fraction $f_t$, the factor $F$ due to electron shielding effect, the pure fast ion current density $J_f$, the net current density $J_{\text{net}} = J_f F$, the electron collision frequency $\nu_e$, the thermal electron bounce frequency $\omega_b$, and the normalized electron collision frequency $\nu_{e\ast}$.

coordinate near the core, where the collisionality $\nu_{e\ast}$ remain small. As we move further outward, the normalized collision frequency $\nu_{e\ast}$ increases rapidly with the radial coordinate. Note that $F$ decreases with the increase in $\nu_{e\ast}$ [11]. The radial increase of $\nu_{e\ast}$ then dominates over the increase in $f_t$, making $F$ decease with the increase in radial coordinate near the edge. As a comparison, we also plot the $F$ for the case of $\nu_{e\ast} = 0$ and the resulting net current. The results show that $F$ for $\nu_{e\ast} = 0$ is always larger than that of $\nu_{e\ast} > 0$ and is monotonically increasing toward the edge.

Next, we evaluate the FLR effects. The FLR effects can appear in the simulations via the following channels: (1) in computing guiding-center position from the particle ionization location, (2) in computing the effective field used in the guiding-center drift, (3) in averaging over Monte-Carlo markers to get fast ion current, and (4) in determining whether a particle touches the first wall. The first channel is always included in the simulation. The (2)-(4) channels are often turned off in order to reduce computational overhead. We compare the results when the (2)-(4) channels are included with those obtained when they are turned off. Fig. 5 shows that the difference between the two cases is negligible, indicating those FLR effects are not important.

Next, we examine the density dependence of the NBCD. We scan the electron density via scaling the profile in Fig. 1 by a factor $\alpha$ with values ranging from $1/32$ to $1.5$.

Fig. 6 plots the radial profiles of fast ion current and net current obtained when different plasma densities are used. The results indicate that both the fast ion current and net current decrease with the increase in plasma density. Another observation is that the driven current density $J_{\text{net}}$ always reaches its peak value at off-axis locations.

Fig. 7 plots the dependence of the volume integrated currents on the density. In the very low density regime ($n_e/n_{\text{ref}} = 1/32$), the shine-through loss is large (about 90%). Even for this case, the NBCD efficiency is larger than all the cases with $n_e/n_{\text{ref}} > 1/32$. As is shown in Fig. 7 the shine-through loss decrease with the density, which is as expected. The shine-through loss is not the dominant factor that determines the drive efficiency (otherwise we
would see the drive efficiency increase with the density. The dominant factor turns out to be the slowing-down time.

With the plasma density increasing, there are two competing factors that influence the NBCD efficiency, namely the edge loss (shine-through loss and orbit loss) and the slowing-down time.

The edge loss depends on the plasma density via several processes: the ionization probability and thus the shine-through fraction and ionization location, the latter of which (weakly) determines the first-orbit loss fraction and determines the ratio between trapped and passing fast ions. With the density increasing, the edge loss usually decreases, which is beneficial for current drive. On the other hand, with the density increasing, the slowing-down time becomes shorter. As a result, fast ions can remain energetic for shorter time and thus less fast ions can contribute to the fast ion current, which is deleterious to the current drive.

These two competing factors can be combined in the following simple way to take into account its effects on the NBCD efficiency:

$$\eta = c \bar{t}_s (1 - L),$$

where $c$ is a constant coefficient, $\bar{t}_s$ is the average slowing down time (averaged over all the fast ions), $L$ is the power loss fraction (due to shine-through and orbit loss). Both $\bar{t}_s$ and $L$ depend on the plasma density and can be directly computed in simulations (collision model used in this work is discussed in Appendix B). Using the computed $\bar{t}_s$ and $L$, we can use Eq. (2) to obtain the dependence of the NBCD efficiency $\eta$ on the plasma density (with a undetermined constant coefficient $c$).

The semi-analytic estimation, Eq. (2), of the dependence of NBCD efficiency on the density is plotted in Fig. 7, which shows good agreement with the results directly obtained in the simulations.

Also shown in Fig. 7 is the heating efficiency, which increases with the increase in plasma density. This increase is mainly correlated with the decrease in the shine-through loss, as is shown in Fig. 7 (b).
4 Summary and discussion

We performed a realistic simulation of neutral beam current drive on the EAST tokamak. We found that the driven current is off-axis although the fast ion current is on-axis. This is due to the electron shielding effect, which is strong near the magnetic axis. We also examine the dependence of the current drive on the plasma density. While the neutral beam heating efficiency increases with the increase in plasma density, the neutral beam current drive efficiency decreases with increase in the density. Therefore a comprise is needed in choosing plasma density in order to optimize both the current drive efficiency and heating efficiency at the same time.

Trapped fast ions carry current via their radial gradient generating bootstrap current. It seems that the electron shielding effect can not shield the bootstrap current in the same way they shield the current carried by passing fast ions. However, in this work, we do not distinguish between the current carried by passing fast ions and that carried by trapped fast ions (bootstrap current), and multiplying the fast ion current by a shielding factor to take into account the electron shielding effect. This may give incorrect results for the bootstrap current. Since the density of fast ions is low and their radial gradient is also low, the bootstrap current should be small. Considering this, the above approximation should give reasonable results.

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5.1 Conflict of interest
The authors have no conflicts to disclose.

6 DATA AVAILABILITY
The data that support the findings of this study are available at https://github.com/Youjunhu/TGCO, which is licensed under the GNU General Public License v3.0.

A Fast ion birth distribution
The neutral beam ionization is modeled by the Monte-Carlo method\[18, 13\]. Typical number of Monte-Carlo markers initially loaded in the simulations is $4 \times 10^4$. Figure 8 plots the two-dimensional distribution of ionized particles in the poloidal plane (averaged over the toroidal direction) and in the toroidal plane (averaged over the vertical direction). The results indicate most fast ions are born on the low-field-side.

It is customary to use $(P_\phi, \Lambda)$ coordinates to describe the guiding-center phase space in order to classify what types of orbits each marker will take in the equilibrium field, where $\Lambda = \mu B_0/E$, $E$ is ion kinetic energy, and $P_\phi$ is the canonical toroidal angular momentum. Figure 9 is this kind of plot for the magnetic configuration specified in Fig. 1.

Fig. 9 shows that non-negligible fraction fast ions are in the loss region (the region between the blue line and red line), which means that the orbits of these particles will touch the LCFS. However, if we use the first wall as the loss boundary (as we do in the simulations), actual loss fraction is much smaller (1%), as is shown in Fig. 10.
A.1 Continuous beam injection

To obtain the steady state of fast ions, we need include the continuous beam source. A straightforward Monte-Carlo implementation of this continuous injection would be to introduce new Monte-Carlo markers to represent the newly injected physical particles at each time step. This method is computationally expensive. For a time-independent background plasma, there is an efficient method that involves only a single injection and then utilizes the time shift invariant to infer the contribution of all the other injections. This method is illustrated in Fig. 11.

The above method works only for a time-independent background plasma and constant beam power. For time-dependent background plasma, re-injecting new Monte-Carlo markers seems to be the only method available. The present work considers a time-independent background plasma with a constant beam power and use the above efficient method.

B Collision model

In the zero drift-orbit width approximation, the time it takes for a fast ion of velocity $v_1$ to be slowed down to $v_2$ by the collision friction with the background ions and electrons is given by

$$\tau_s = \frac{\tau_s}{3} \ln \left[ \frac{1 + \left( \frac{v_2}{v_1} \right)^3}{\left( \frac{v_2}{v_1} \right)^3 + \left( \frac{v_2}{v_1} \right)^3} \right],$$

(3)
Figure 9: NBI fast ion markers (blue points) mapped to the \((P_\phi, \Lambda)\) plane (only ten percent of total markers are shown here). Also shown are the passing trapped boundary, the magnetic axis, high-field-side and low-field side of last-closed-flux-surface (LCFS) for fast ions of \(E = 55\) keV in the magnetic configuration specified in Fig. 1. Here \(B_n = 1T\) and \(L_n = 1m\).

where \(v_c\) is the critical velocity defined by

\[
v_c = \left( \frac{3\sqrt{\pi} m_e}{4 n_e} \sum_i \frac{n_i Z_i^2}{m_i} \right)^{1/3} \sqrt{\frac{2T_e}{m_e}},
\]

(4)

\(\tau_s = 1/\nu_s\) with \(\nu_s\) defined by

\[
\nu_s = \frac{4 m_f}{3\sqrt{\pi} m_e} \frac{\Gamma_{f/e}}{m_e (2T_e/m_e)^3/2},
\]

(5)

which is the slowing-down rate due to background electrons,

\[
\Gamma_{f/e} = \frac{n_e Z_f^2 e^4}{4\pi \nu_d^2 m_f^2} \ln \Lambda,
\]

(6)

\(\ln \Lambda\) is the Coulomb logarithm (\(\ln \Lambda = 24 - \ln \left(10^{-6} \sqrt{n_e/T_e}\right) \approx 16.97\) is used in this work, where \(n_e\) and \(T_e\) are in units of \(m^{-3}\) and keV, respectively).

Both \(\tau_{se}\) and \(v_c\) have a radial dependence via their dependence on the plasma density and temperature. The radial profile of \(\tau_s\) for the plasma specified in Fig. 1 is plotted in Fig. 12 for fast ions of 55 keV to be slowed down to a cutoff velocity (chosen as \(mv^2/2 = 2T_i(0)\) in this article). The result shows that typical slowing-down time in the core region is about 40 ms.

The formula (5) only includes the slowing-down, neglecting pitch-angle scattering and energy diffusion. For realistic simulations that include pitch-angle scattering, energy diffusion, finite-orbit-width effect, and non-uniform plasma profiles, we need to use numerical integration to determined the slowing-down process of fast ions. The Monte-Carlo algorithm used in this work for collision of fast ions with background electrons and ions is specified in Ref[20], where the pitch-angle variable \(\lambda = v_\parallel/v\) and velocity \(v\) are altered at the end of each time step according to the following scheme:

\[
\lambda_{\text{new}} = \lambda(1 - \nu_d \Delta t) \pm \sqrt{(1 - \lambda^2)\nu_d \Delta t},
\]

(7)
Figure 10: Time evolution of the loss fraction and thermalization fraction for a pulse NBI fast ions born at \( t = 0 \). The loss fraction during the 100 ms is about 1%.

and

\[
v_{\text{new}} = v - \nu_d \Delta t \left[ \frac{1}{3} + \frac{v^3}{v^3} \right] + \frac{\nu_s \Delta t}{m_f v} \left[ T_e - \frac{1}{2} T_i \left( \frac{v_e}{v} \right)^3 \right] \pm \sqrt{\frac{\nu_s \Delta t}{m_f} \left[ T_e + T_i \left( \frac{v_e}{v} \right)^3 \right]},
\]

(8)

where the blue term corresponds to the energy diffusion, \( \pm \) denotes a randomly chosen sign with equal probability for plus and minus, \( \Delta t \) is the time step, \( \nu_d \) is the pitch-angle scattering rate given by

\[
\nu_d = \left( \frac{Z_{\text{eff}}}{v^3} \right) n_e Z_f^2 e^4 \ln \Lambda_e,
\]

(9)

where \( Z_f \) is the charge number of fast ions, \( \nu_s \) is the slowing down rate given by Eq. (5).

The steady-state profiles of fast ion pressure, ion and electron heating power density for the reference case are plotted in Fig. 13.

C Formulas for electron shielding effect

The electron collision time (characterizing electron collision with ions) is [19]

\[
\tau_e = \frac{12\pi^{3/2}}{e_0} \frac{e^2}{\sqrt{m_e}} \frac{\varepsilon^3}{T_e^{3/2}} \ln \Lambda_e,
\]

(10)

where \( \ln \Lambda_e \) is the Coulomb logarithm for electron-ion collision (\( \ln \Lambda_e = 15.2 - 0.5 \ln(n_e/10^{20}) + \ln(T_e) \) is used in this work, where \( n_e \) and \( T_i \) are in units of \( m^{-3} \) and keV, respectively). The dimensionless electron collision parameter \( \nu_{e*} \) is defined by

\[
\nu_{e*} = \frac{\nu_e}{\varepsilon \omega_{pe}},
\]

(11)

where \( \nu_e = 1/\tau_e \), \( \varepsilon = r/R_0 \) is the local inverse aspect ratio, \( r = (R_{\text{max}} - R_{\text{min}})/2 \), \( R_0 = (R_{\text{min}} + R_{\text{max}})/2 \), \( R_{\text{min}} \) and \( R_{\text{max}} \) are, respectively, the maximal and minimal values of the
Figure 11: An efficient way of simulating multiple injections. The contribution of injections at $t = j\Delta t$ with $j = 0, 1, 2$ to the fast ion distribution at $t = 3\Delta t$ can be obtained by following the time history of the particles injected at $t = 0$ and recording their contribution to the fast ion distribution at $t = \Delta t, 2\Delta t, 3\Delta t$, indicated respectively by $A_1$, $A_2$, and $A_3$. Then it is ready to see that $B_1 = A_1$, $B_2 = A_2$, and $B_3 = A_3$. Therefore the contributions of the multiply injections can be inferred from the time history of a single injection.

cylindrical coordinate $R$ on a magnetic surface, $\omega_{be}$ is the typical bounce (angular) frequency of thermal electrons, which is given by (in the approximation of zero-orbit-width and deeply trapped electrons)

$$\omega_{be} = \frac{\sqrt{2T_e/m_e}}{qR_0} \left(\frac{\varepsilon}{2}\right)^{1/2},$$

(12)

where $q$ is the safety factor of the magnetic surface. Using this, Eq. (11) is written as

$$\nu_{e*} \equiv \frac{\nu_e}{\varepsilon^{3/2} \sqrt{T_e/m_e}/(qR_0)}.$$

(13)

Due to the electron shielding effect, the ratio of the net current to the fast ion current is given by

$$F \equiv \frac{\langle j_B \rangle}{\langle j_f \parallel B \rangle} = 1 - \frac{Z_f}{Z_{eff}} (1 - L_{31}),$$

(14)

where $\langle \ldots \rangle$ is the magnetic surface averaging, $L_{31}$ is the bootstrap current coefficient before the electron density gradient. The formula of $L_{31}$ given by Sauter et al. [10] is

$$L_{31} = \left(1 + \frac{1.4}{Z_{eff} + 1}\right) X - \frac{1.9}{Z_{eff} + 1} X^2 + \frac{0.3}{Z_{eff} + 1} X^3 + \frac{0.2}{Z_{eff} + 1} X^4,$$

(15)

with

$$X = \frac{f_t}{1 + (1 - 0.1f_t)\sqrt{\nu_e^*} + 0.5(1 - f_t)\nu_{e*}/Z_{eff}},$$

(16)

$f_t$ is the effective trapped fraction [21, 22].

$$f_t = 1 - \frac{3}{4} \left(\frac{B^2}{B_{max}^2}\right) \int_0^1 \frac{\lambda d\lambda}{\sqrt{1 - \lambda B/B_{max}}},$$

(17)

where $B_{max}$ is the maximal value of magnetic field strength on a magnetic surface.
Figure 12: Radial profiles of slowing-down time for fast ions of kinetic energy 55 keV (assuming zero drift-orbit width) to slow down to the cutoff energy $2T_\text{i}(0)$ in the equilibrium specified in Fig. 1.

Figure 13: Steady-state profiles of fast ion pressure, ion and electron heating power density. The ion heating power is 0.426 MW, and electron heating power is 0.443 MW.

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