

Folded double layer: a novel nonlinear ion acoustic wave propagating obliquely in a magnetized plasma

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Abstract

A novel nonlinear structure is reported with a ‘near-unipolar’ solution which is obtained analytically by using Sagdeev pseudopotential technique for a three component magnetized plasma. The plasma is consisting of warm fluid ions and two types of electrons having Maxwell Boltzmann distributions. The structure is an amalgamation of the newly discovered super solitary wave and a conventional double layer. Because of the extra ‘wiggle’ or ‘fold’ in the usual monopolar electric field, it is named as ‘Folded Double Layer’. The presence of the extra wiggle, or ‘fold’, in the otherwise monopolar pulse may well modify the particle acceleration processes which remains hitherto unknown, or unnoticed, during the satellite observations.

Keywords: nonlinear waves, double layers, magnetized plasma

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

1. Introduction

The boundary layers of the Earth’s magnetosphere are the hotbed of conventional and non-conventional coherent nonlinear structures in the electric field (E-field) data which are formed due to the large discontinuities in the background plasma conditions. The theory of the observed bipolar and monopolar pulses in the electrostatic solitary waves [1] are often implicitly related to the mathematical formalism of nonlinear coherent structures like solitons [2] and kink solitons [3]. However, the essential condition of integrability for the latter ones deters their direct applicability in the space plasma. The Sagdeev pseudopotential technique [4] often provides an useful bridge between the two. Given the background plasma condition, it allows one to predict different kinds of nonlinear coherent structures without addressing the problem of integrability altogether, though it retains the physical characteristics and boundary conditions of their mathematical counterparts [5]. In recent days, extra-nonlinear structures like super solitary waves (SSWs) [6–9] and flat top solitary waves [10], which have been obtained theoretically by using Sagdeev pseudopotential technique, are finding their ways in interpreting more composite and non-conventional

pulses recorded in the E-field data during the satellite observation [11]. With all their finer differences in their morphologies, they may together be called by an umbrella name of extra nonlinear solitary waves (ENLSWs), or, more generally, as modified solitary waves, since they all ascertain a ‘no net potential drop (or rise)’ across the localized pulse. In other words, they all satisfy the recurrence condition of a conventional soliton. On the other hand, a monopole is necessarily associated with a net change in the potential across the pulse [12]. Double layer (DL) structures, observed as monopolar pulses in the E-field data during satellite expeditions, play an important role in the particle acceleration and transport in the boundary layers of the Earth’s magnetosphere [12, 13]. This poses an open question whether a monopolar or DL-like structure, too, can have a non-conventional morphology. The question is important since such a non conventionality would surely affect the particle acceleration process at the magnetospheric boundary layer. Here, for the first time, we are theoretically predicting a non-conventional ion acoustic double layer, propagating obliquely in a magnetized plasma, which appears with an extra wiggle in their otherwise monopolar pulses. Due to their extra ‘fold’, we call them ‘Folded Double Layers’ (FDLs). They trivially determine the

termination of the associated SSWs in the parameter space. An SSW was called ‘super’ because of their extra large amplitude and speed compared to their associated DL [14]. Here, since their ‘superiority’ (larger amplitude) over the DL is lost, we propose to call them as Folded Solitary Waves (FSWs). In the present paper, we shall use this latter nomenclature synonymously with SSW. These FSWs are 1D steady state solitary structures which are essentially different from ‘foldons’ in $2 + 1$ dimensions [15].

The paper is organized as follows: the following section (section 2) provides the analytical solution, section 3 gives the numerical results and discussions and finally section 4 gives the concluding remarks of the paper.

2. The analytical solution

We have considered a three component plasma where the warm fluid ions are magnetized and the two types of electrons having Boltzmann distributions are separately in thermal equilibrium. The ambient magnetic field B_0 is lying in the z direction while the wave is propagating in the $y - z$ plane making an angle θ with the magnetic field.

The normalized fluid equations are as follows:

$$\frac{\partial n_i}{\partial t} + \nabla \cdot (n_i \mathbf{V}_i) = 0, \quad (1)$$

$$\frac{\partial \mathbf{V}_i}{\partial t} + (\mathbf{V}_i \cdot \nabla) \mathbf{V}_i = -\nabla \Phi + \alpha_i (\mathbf{V}_i \times \hat{\mathbf{b}}) - 3\sigma_i n_i \nabla n_i. \quad (2)$$

The Boltzmann distribution of electrons are given by

$$n_e = \mu \exp\left(\frac{\Phi}{\mu + \nu\beta}\right) + \nu \exp\left(\frac{\beta\Phi}{\mu + \nu\beta}\right), \quad (3)$$

where i (e) represents ions (electrons), $\alpha_i = \frac{\Omega_i}{\omega_{pi}}$ is the ratio of the ion cyclotron frequency to the ion plasma frequency, $\hat{\mathbf{b}}$ is the unit vector along the magnetic field, $\sigma_i = \frac{T_i}{T_{\text{eff}}}$ represents the ratio of the ion temperature to the effective temperature of the electrons, $T_{\text{eff}} = \frac{T_{ec}T_{eh}}{\mu T_{eh} + \nu T_{ec}}$ being the effective electron temperature, μ (ν) are the ambient densities of cooler (hotter) electrons and $\beta = \frac{T_{ec}}{T_{eh}}$ is the temperature ratio of cooler to hotter electrons.

We have normalized all the densities by the ambient plasma density n_0 , time by the inverse of the ion plasma frequency ω_{pi}^{-1} , length by the effective Debye length, pressure by the ion equilibrium pressure $P_0 = n_0 T_i$, temperature by the effective electron temperature T_{eff} , magnetic field by the ambient magnetic field B_0 and the potential is normalized by $\frac{T_{\text{eff}}}{e}$. All the velocities are normalized by the effective ion acoustic speed $C_{\text{eff}} \left(= \sqrt{\frac{T_{\text{eff}}}{m_i}} \right)$. Ideally, the Mach number is defined as the wave velocity normalized by the linear acoustic speed of the corresponding model, obtained from its dispersion relation [16]. For a magnetized plasma, however, this would make our analytical formalism more complicated. For the sake of convenience, instead we choose C_{eff} as the normalizing acoustic speed for the wave velocity. The

corresponding Mach number is defined as the effective Mach number M_{eff} [7, 8, 17].

We have considered the following boundary conditions

$$\begin{aligned} \text{at } |\mathbf{r}| \rightarrow \infty, \quad \Phi &\rightarrow 0, \\ V_i &\rightarrow 0, \quad n_i \rightarrow 1, \quad \text{and} \quad P_i \rightarrow n_0 T_i, \end{aligned} \quad (4)$$

where $\mathbf{r} = \mathbf{k}_y y + \mathbf{k}_z z$ is the corresponding position vector, $\mathbf{k}_y, \mathbf{k}_z$ being the direction cosines along the y and z directions respectively.

In order to have the stationary state solution, we introduce the generalized coordinate

$$\eta = k_y y + k_z z - M_{\text{eff}} t. \quad (5)$$

It is customary for an unmagnetized plasma to solve the density of the each fluid species in the stationary state in terms of the potential Φ and then to solve the Poisson’s equation implementing the Sagdeev pseudopotential

$$\frac{d^2 \Phi}{d\eta^2} = n_i - n_e = -\frac{\partial \Psi_{\text{un}}(\Phi)}{\partial \Phi} \quad (6)$$

which leads to

$$\frac{1}{2} \left(\frac{d\Phi}{d\eta} \right)^2 + \Psi_{\text{un}}(\Phi) = 0, \quad (7)$$

where Ψ_{un} is defined as the Sagdeev pseudopotential and the subscript ‘un’ indicates the unmagnetized plasma.

In the case of a magnetized plasma, however, only the density of an unmagnetized species can be expressed in terms of Φ explicitly. The density of a single magnetized species, however, can only be estimated implicitly by assuming the charge neutrality condition [18]

$$n_i \approx n_e. \quad (8)$$

So, we have incorporated the charge neutrality condition (equation (8)) instead of the Poisson’s equation for our analysis.

After transforming the space and time derivatives with the generalized coordinate η and applying appropriate boundary conditions (equation (4)), the continuity equation (equation (1)) gives us

$$n_i = \frac{M_{\text{eff}}}{M_{\text{eff}} - k_y V_{iy} - k_z V_{iz}}. \quad (9)$$

Using equation (9) and systematically eliminating V_{ix}, V_{iy} and V_{iz} from all the three components of the momentum equation (equation (2)) [19], we get

$$\begin{aligned} \frac{d}{d\eta} \left[\frac{1}{2} \frac{1}{n_i} \frac{d^2}{d\eta^2} \left(\left(\frac{M_{\text{eff}}}{n_i} \right)^2 + 3\sigma_i n_i^2 + 2\Phi \right) \right. \\ \left. + \alpha_i^2 \left(\frac{1}{n_i} + \frac{k_z^2}{M_{\text{eff}}^2} (\sigma_i n_i^3 + \tilde{n}_i) \right) \right] = 0, \end{aligned} \quad (10)$$

$$\tilde{n}_i = \int n_i d\Phi. \quad (11)$$

Integrating equation (10) and using boundary conditions in equation (4), we get the second order ordinary differential equation

$$\frac{d^2 f}{d\eta^2} + \alpha_i^2 g = 0, \quad (12)$$

where

$$f = \frac{1}{2} \left(\frac{M_{\text{eff}}^2}{n_i^2} + 3\sigma_i n_i^2 \right) + \Phi, \quad (13)$$

and

$$g = 1 + \frac{k_z^2}{M_{\text{eff}}^2} (\sigma_i n_i^4 + n_i \tilde{n}_i) - n_i \left(1 + \frac{k_z^2}{M_{\text{eff}}^2} (\sigma_i + \delta) \right) \quad (14)$$

with

$$\delta = \tilde{n}_i|_{\Phi=0}. \quad (15)$$

Integrating once again, equation (12) reduces to

$$\frac{1}{2} F^2 \left(\frac{d\Phi}{d\eta} \right)^2 + \alpha_i^2 \int (gF) d\Phi = 0, \quad (16)$$

where

$$F(\Phi) = \frac{df}{d\Phi} = 1 + \left(3\sigma n_i - \frac{M_{\text{eff}}^2}{n_i^3} \right) \frac{dn_i}{d\Phi}. \quad (17)$$

This gives us the analog of the familiar ‘energy equation’ (equation (7)) of the Sagdeev pseudopotential for magnetized plasma

$$\frac{1}{2} \left(\frac{d\Phi}{d\eta} \right)^2 + \Psi(\Phi) = 0, \quad (18)$$

where $\Psi(\Phi)$ is the Sagdeev pseudopotential

$$\Psi(\Phi) = \frac{\alpha_i^2 L_i}{F^2}, \quad (19)$$

L_i and F being the algebraic functions of all the parameters.

The numerator of the Sagdeev pseudopotential thus becomes

$$\begin{aligned} L_i(\Phi) &= \int (gF(\phi)) d\phi \\ &= \left[k_y^2 \Phi + \frac{1}{2} M_{\text{eff}}^2 \left(\frac{S}{n_i} \right)^2 + w - 3h \right] \\ &\quad + k_z^2 \left[h - \frac{1}{n_i} (\sigma S + \tilde{S}) + \frac{1}{2} \left(\frac{w}{M_{\text{eff}}} \right)^2 \right], \end{aligned} \quad (20)$$

where

$$S = 1 - n_i; \quad \tilde{S} = \delta - \tilde{n}_i, \quad (21)$$

$$h = \frac{\sigma(1 - n_i^2)}{2}; \quad t = \sigma(1 - n_i^3) \quad \text{and} \quad w = \tilde{S} + t \quad (22)$$

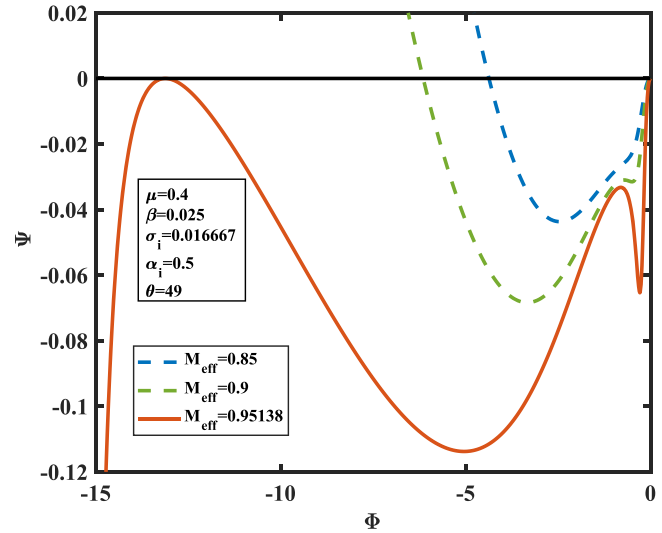


Figure 1. Sagdeev pseudopotential of FDL (solid curve) and FSWs (dashed curves).

F^2 becomes the denominator of the Sagdeev pseudopotential.

The boundary condition at the equilibrium ($\Phi = 0$) is given as

$$\Psi(\Phi)|_{\Phi=0} = \frac{\partial \Psi(\Phi)}{\partial \Phi} \Big|_{\Phi=0} = 0 \quad \text{and} \quad \frac{\partial^2 \Psi(\Phi)}{\partial \Phi^2} \Big|_{\Phi=0} < 0. \quad (23)$$

The last condition ascertain the negative value of the pseudopotential, i.e.

$$\Psi(\Phi) < 0 \quad \text{for} \quad \begin{cases} 0 < \Phi < \Phi_0 & \text{when } \Phi_0 > 0 \\ \Phi_0 < \Phi < 0 & \text{when } \Phi_0 < 0 \end{cases}. \quad (24)$$

The Sagdeev pseudopotential will ascertain the existence of a nonlinear structure provided it has a second root at some $\Phi = \Phi_0$ other than zero. The slope of Ψ at Φ_0 determines the type of the solution, being $\frac{\partial \Psi}{\partial \Phi} \Big|_{\Phi=\Phi_0} \neq 0$ for a solitary wave

and $\frac{\partial \Psi}{\partial \Phi} \Big|_{\Phi=\Phi_0} = 0$ for a DL.

3. Numerical results

To facilitate our model, we have chosen a convenient set of plasma parameters having ambient densities of the cooler electrons $\mu = 0.4$, the cooler to hotter electron temperature ratio $\beta = 1/40$, the ratio of the ion temperature to the effective electron temperatures $\sigma_i = 1/60$ while $\alpha_i = 0.5$ is the ratio of the ion cyclotron and ion plasma frequency. The obliqueness is considered as $\theta = 49^\circ$ and the effective Mach numbers are considered as $M_{\text{eff}} = 0.85, 0.9$ and 0.951385 .

To analyze the trajectory of the pseudoparticle, in figure 1 we have plotted the Sagdeev pseudopotential for all the above mentioned parameters while figure 2 shows the phase portraits of the corresponding pseudoparticle.

For $M_{\text{eff}} = 0.95138$, the solid curve in figure 1 satisfies the usual boundary condition for a DL [20] ($\frac{\partial \Psi}{\partial \Phi} \Big|_{\Phi=\Phi_0} = 0$) at

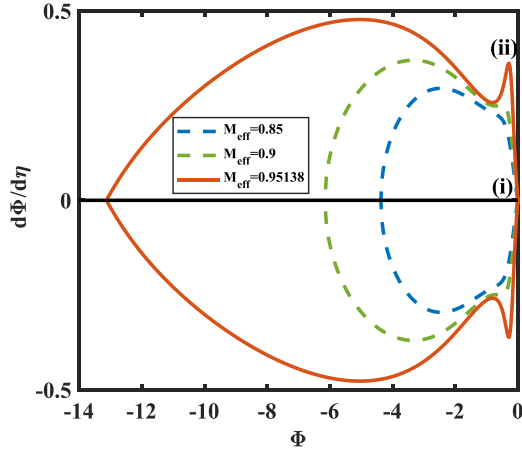


Figure 2. Phase portrait of FDL (solid curve) and FSWs (dashed curves).

its posterior part while its anterior part reveals the spiky subwell of an FSW (or SSW) in a magnetized plasma. The amalgamation of these two signatures provides this unique structure which we choose to call a ‘FDL’. This particular solution is preceded by a series of FSWs (dashed lines), or more popularly known as SSWs, which show their typically spiky ‘auxiliary’ subwells at the lower potential [7, 8]. Clearly, the observed spiky ‘fold’ in the FDL (solid curve) is the reminiscent of the spiky subwells of the preceding FSWs. We have previously shown that, for a magnetized plasma, the subwells of FSWs arise due to the imminent singularity in the Sagdeev pseudopotential [8]. It appears to be true for the extra ‘fold’ or subwell of the FDL as well.

The FDL solution in figure 1 is marked by a super large amplitude. The novelty of the FDL becomes further apparent in the following phase portrait (figure 2). The separatrix of the FDL engulfs all the extra-nonlinear FSWs within it in exactly the same way as a conventional DL is known to terminate a conventional ‘rarefactive’ (potential well, negative amplitude in this case) solitary wave. The extra large amplitude of the FDL is worth to note. Certainly, it is far less likely to occur for a ‘bounded’, or ‘potential hill’ type of solution. For an ‘unbound’, or ‘potential well’ type structures in an unmagnetized plasma, the increasing effective Mach number/amplitude often found to smoothen the ‘fold’ of the FSW, thus reducing, but certainly not eradicating, the possibility of an FDL. Conversely, for a magnetized plasma, the depth of the auxiliary subwell causing the ‘fold’ increases with the increasing M_{eff} and amplitude. With a faster rate of increase in the fold with M_{eff} compared to its amplitude, the solution would hit the ‘point of singularity’ in the parameter space prior to the formation of any FDL and any physical solution for FSW will cease to exist. For the present set of parameters, the trend is just the opposite. The ‘point of singularity’ lies well beyond the regime which supports any kinds of physical solution, allowing an FDL to form. The occurrence of FDL is, however, conditional and may be qualitatively expressed by

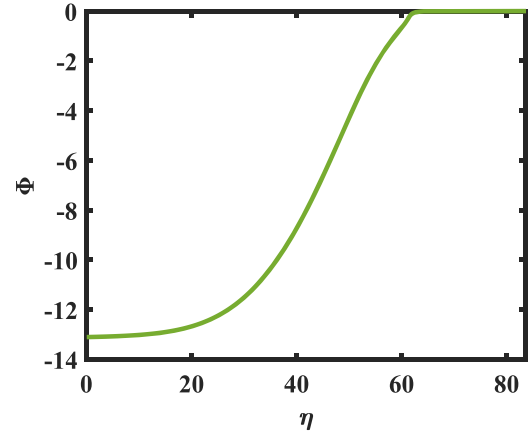


Figure 3. Potential profile of FDL.

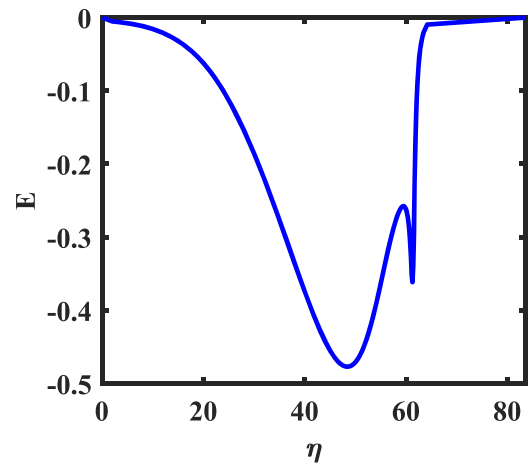


Figure 4. Electric field profile of FDL.

the probability factor p , defined as

$$p = \frac{d\Phi_0/dM}{dl/dM}, \quad (25)$$

where l is the total length of the fold between the points (i) and (ii) marked in figure 2. We conjecture that, to have a FDL, the condition $p > 1$ needs to be satisfied. The very large amplitude of the solution is further get highlighted in figure 3 which represents the potential profile of the FDL, associated with the solid curve in figure 1, while figure 4 shows the corresponding electric field profile. The prominent wiggle in the electric field profile readily distinguishes it from the usual monopolar pulses observed in the space. We may also categorize this FDL as an extra nonlinear double layer (ENDL) in comparison with the previously mentioned ENLSWs.

Before proceeding further, let us recapitulate the different nomenclatures associated with the umbrella name of ENLSW. They are namely the generalized variable solitary waves (gVSWs), the curve of inflection (CoI) and the fully grown SSW or FSW. A CoI may again be of two types, positive (p-CoI) and negative (n-CoI) [21]. The Sagdeev pseudopotentials of all those structures satisfy the boundary condition of the solitary wave, thus ensuring the recurrence of the initial state, but their morphology differs from that of the regular

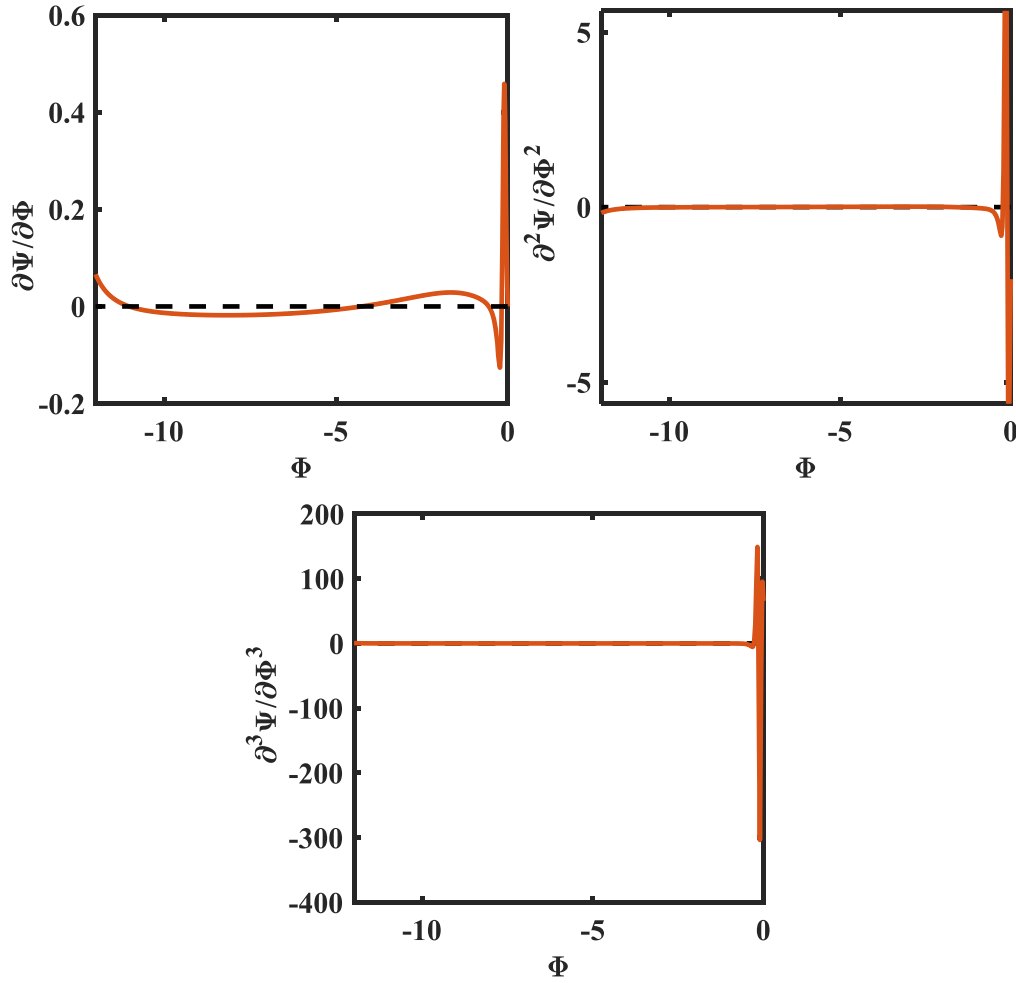


Figure 5. The derivatives of Sagdeev pseudopotential for FDL.

one, the wiggle being most prominent for the FSW and least for the gVSW. Following these properties, they were defined according to the number of roots of the different orders of derivatives of the Sagdeev pseudopotential [22]. The derivative profiles of Sagdeev pseudopotential for FDL is given in figure 5. For better clarity the zoomed version is given in figure 6. A complete list of those conditions are given in the table 1.

To understand the condition (equation (25)) more clearly, we have delineated the parameter space in detail. It reveals that the occurrence of the FDL is highly sensitive to the particular combination of the parameters and even only a slight deviation in any of them may disrupt the solution altogether. Figure 7 shows four snapshots of the Sagdeev pseudopotential which represent the DL and DL-like terminating solutions. We have chosen a convenient parameter set with $\mu_i = 0.35$, $\sigma_i = 1/60$, $\alpha_i = 0.5$ and $\theta = 49^\circ$ while we changed the β value for each set of Ψ .

A careful consideration reveals two major zones in the parameter space, namely

$$\text{Zone A: RSW} \rightarrow \text{RDL}, \quad (26)$$

$$\begin{aligned} \text{Zone B: RSW} &\rightarrow \text{ENLSW} \\ &\rightarrow \text{ENDL} \rightarrow \text{leading to singularity}, \end{aligned} \quad (27)$$

where RSW (RDL) denote regular solitary wave (regular double layer) and ENLSW (ENDL) denote ENLSW (ENDL) respectively. Expectedly, zone A (equation (26)) covers a wide region of the parameter space where conventional or regular rarefactive solitary waves terminates to the corresponding double layers, a well known phenomena for the negative amplitude ion acoustic solitary waves.

It also readily reveals that all the other 3 figures, viz., figures 7(b)–(d) belong to zone B (equation (27)). We note that, in figure 7(d), the solution terminates as a solitary wave (FSW) rather than a DL. The overall transition is

$$\begin{aligned} \text{Type B1: RSW} &\rightarrow \text{gVSW} \rightarrow \text{p-CoI} \\ &\rightarrow \text{FSW} \rightarrow \text{Singularity}. \end{aligned} \quad (28)$$

We may well call it a transition of type B1 (equation (28)). It is identical to that presented in our previous papers [7, 8], and is closely associated with that of Steffy and Ghosh, 2017 [21]. Clearly, the condition of equation (25) is not satisfied here and hence there is no DL solution. Instead, with increasing M_{eff} , the depth of the auxiliary well increases rapidly and indefinitely, leading to the ‘point of singularity’ in the Sagdeev pseudopotential curve where the denominator of equation (17) tends to zero. The associated β value is $1/50$.

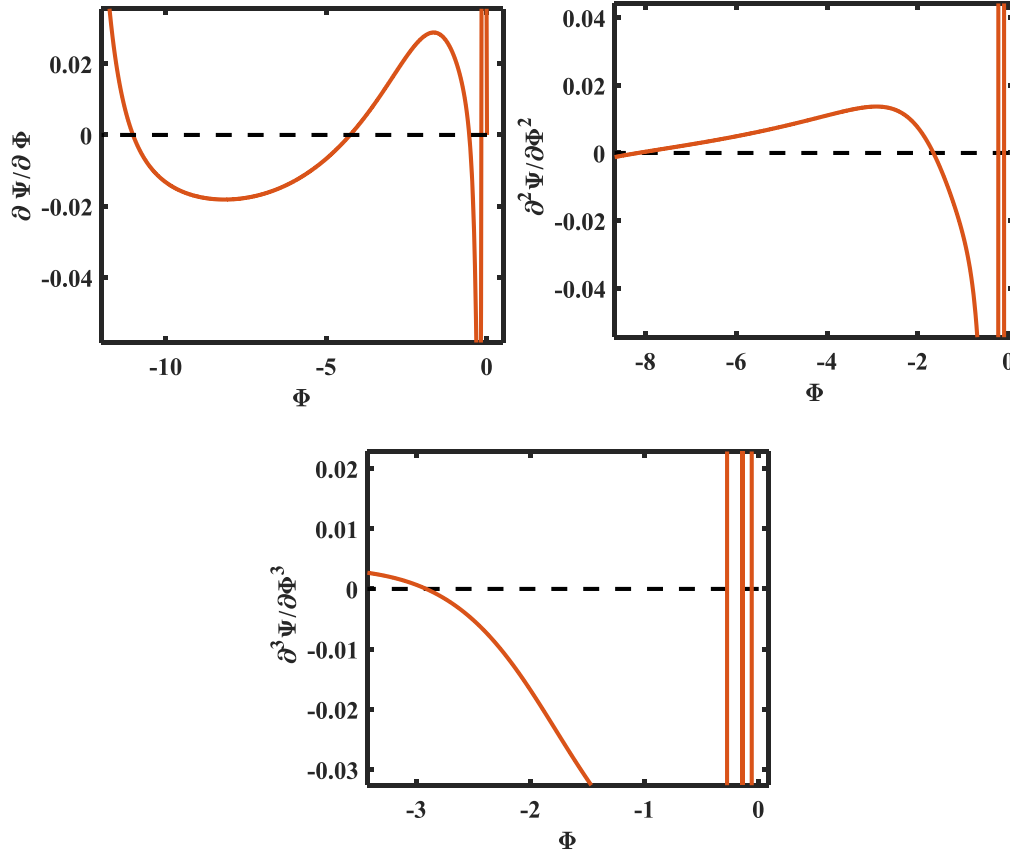


Figure 6. The zoomed version of derivatives of Sagdeev pseudopotential for FDL.

Table 1. Number of roots for the derivatives of Sagdeev pseudopotential for different structures.

Structure	1st derivative	2nd derivative	3rd derivative
RSW	2	1	1
FSW	4	3	3
RDL	3	2	2
VDL	3	4	4
FDL	5	4	4

As we increase the β value from $1/50$ to $1/40$ (figure 7(c)), there appears a fully grown FDL prior to the Sagdeev pseudopotential curve going to the ‘point of singularity’. Beyond FDL, all physical solution ceases to exist. It also ensures the condition of equation (25). Characteristically, the solid curve is the same as that presented in figure 1. The overall transition is now named as type B2 and can be presented as

$$\text{Type B2: RSW} \rightarrow \text{gVSW} \rightarrow \text{p-CoI} \rightarrow \text{FSW} \rightarrow \text{FDL} \rightarrow \text{Singularity.} \quad (29)$$

The Sagdeev pseudopotential curves beyond the FDL would show singularities though they would no longer be physically relevant for our solutions.

As β increases further, i.e. $\beta = 1/30$ in figure 7(b), the wiggle of the FDL gets smoothened. Though it loses any prominent auxiliary subwell, the morphology does not

resemble an RDL either. The terminating DL solution has a slight bump which is analogous to a gVSW. Hence we choose to call them a variable double layer or VDL. The overall transition process is called B3 represented by equation (30)

$$\text{Type B3: RSW} \rightarrow \text{gVSW} \rightarrow \text{VDL} \rightarrow \text{Singularity.} \quad (30)$$

Both FDL and VDL satisfy the boundary condition of RDL but their morphologies differ. We may well club the first two together to call them ENDL. Hence transition B2 (equation (29)) and B3 (equation (30)), together, represent zone B (equation (27)) and B1 (equation (28)) may be considered as a special case of Zone B (equation (27)) which misses the terminating DL.

Though for each case we assumed the initial solution to be an RSW, in zone B (equation (27)), we could not detect any RSW for our chosen parameter set. Like any ENLSW, both FDL and VDL can be accurately defined by the number of roots for the derivatives of the corresponding Sagdeev pseudopotential which have been listed in the table 1.

The four snapshots, together, show that, with increasing β , the solution proceeds from zone B (equation (27)) to zone A (equation (26)), i.e. the extra nonlinear features of the solution disappears turning it to a regular one. Apart from that, within zone B (equation (27)), increase in β pushes the imminent singularity further away allowing an FDL to form. Further increase in β smoothen the wiggle turning a FDL to a VDL and then to a RDL.

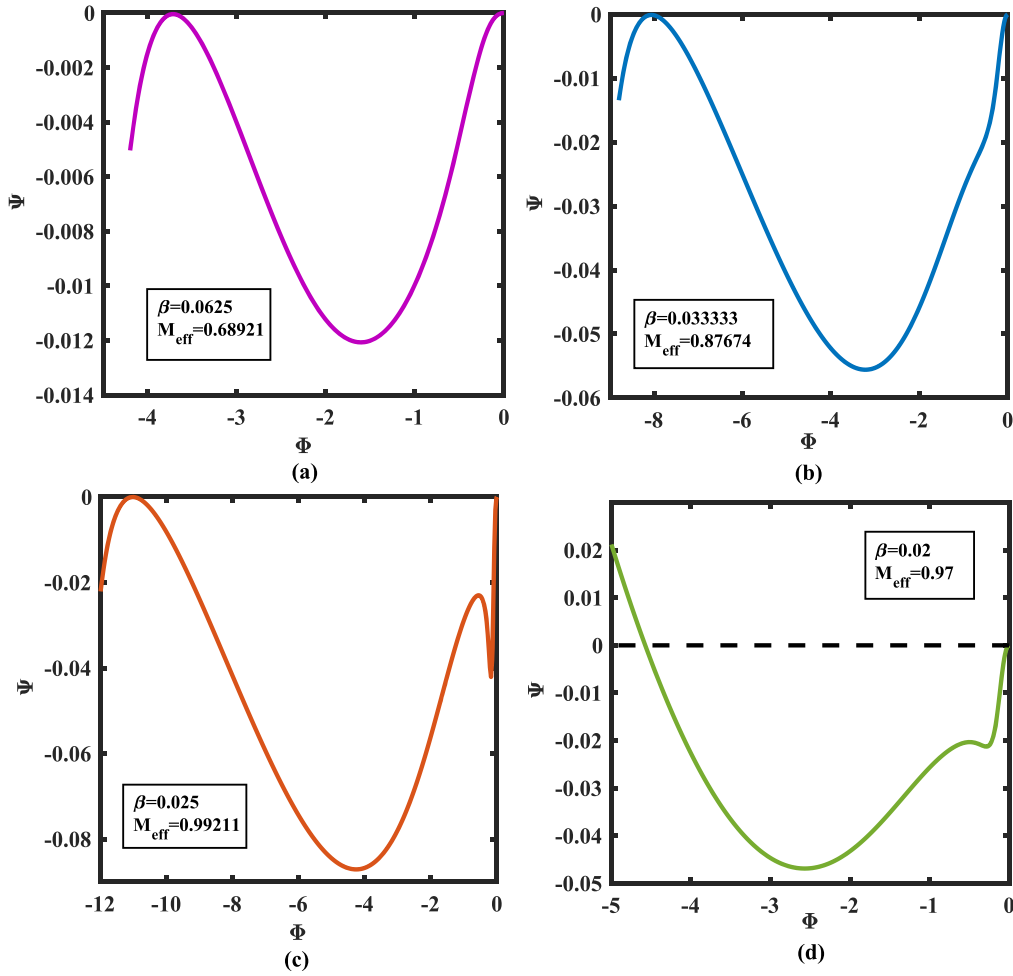


Figure 7. Sagdeev pseudopotential of various structures: (a) DL (b) VDL (c) FDL and (d) FSW.

So far it is found that the FDL occurs for some selected sets of parameters rather than any range of it. Two competing processes, viz., transformation of a rarefactive SW to a DL and appearance of a ‘fold’ due to the imminent singularity result into a FDL according to the condition in equation (25). With such fine balance of parameters, the occurrence of an FDL in reality, is expected to be a rare event which will be hard to detect. Until very recently, there was no recorded events of any FDL or similar kinds of structures in the literature. Moreover, while large amplitude bipolar pulses are fairly well known, recorded events of large amplitude monopolar pulses are not known so far. This raises the question of the feasibility of finding such solutions in reality. Interestingly, large amplitude ‘near unipolar’ structures have been reported very recently in the plasma sheet boundary layers from magnetospheric multiscale spacecraft observations [23]. One of their observed structures actually shows the presence of ‘near unipolar’ (double layer) structure which appears to indicate the presence of a heteroclinic counterpart of FSW. The morphology of one of this structures apparently resembles the FDL solution reported in the present work. It is also observed that such large amplitude ‘near unipolar’ events in E_{\parallel} occur at the region of magnetic reconnection and are expected to be associated with the turbulence and dissipations

at that region. Though they are expected to be a rare event, nevertheless, it shows that, with careful observation and modeling, there is a fair possibility of detecting a FDL like structure in the Earth’s magnetosphere. They may well unfold the intricate micro-physics of the particle transport and acceleration processes due to their unique spiky structures.

4. Conclusions

To summarize, we have obtained the ion acoustic ‘FDL’ for a three component magnetized plasma comprising warm fluid ions and two temperature electrons having Boltzmann distributions. The novel structure is an extra-nonlinear DL where, for the first time, a monopolar pulse has appeared with an extra ‘fold’. The theoretical prediction may well find its way into the satellite observations which were possibly remained unnoticed so far. The structure is preceded by a series of ‘wiggled’ solitary waves which we proposed to call as the ‘FSW’, rather than a ‘super’ one. The lateral inversion of the auxiliary subwell to the lower potential, and the imminent singularity, the two hallmarks for the FSW in a magnetized plasma, are appeared to be responsible for the

emergence of the FDL, making it less likely for an unmagnetized plasma.

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