Published paper

Istomin, Y., Pokhotelov, O.A., Balikhin, M.A. (2009) *Nonzero electron temperature effects in nonlinear mirror modes*, Physics of Plasmas, 16 (12), Art No.122901
http://dx.doi.org/10.1063/1.3275787

Nonzero electron temperature effects in nonlinear mirror modes

Ya. N. Istomin*

Lebedev Physical Institute, Leninsky Prosp., 53, Moscow, 119991, Russia

O. A. Pokhotelov† and M. A. Balikhin‡

Automatic Control and Systems Engineering,

University of Sheffield, Sheffield S1 3JD, United Kingdom

(Dated:)

Abstract

The nonlinear theory of the magnetic mirror instability (MI) accounting for nonzero electron temperature effects is developed. Based on our previous low-frequency approach to the analysis of this instability and including nonzero electron temperature effects a set of equations describing nonlinear dynamics of mirror modes is derived. In the linear limit a Fourier transform of these equations recovers the linear MI growth rate in which the finite ion Larmor radius and nonzero electron temperature effects are taken into account. When the electron temperature T_e becomes of the same order as the parallel ion temperature $T_{\|}$ the growth rate of the mirror instability is reduced by the presence of a parallel electric field. The latter arises because the electrons are dragged by nonresonant ions which are mirror accelerated from regions of high to low parallel magnetic flux. The nonzero electron temperature effect also substantially modifies the mirror mode nonlinear dynamics. When $T_e \simeq T_{\|}$, the transition from the linear to nonlinear regime occurred for wave amplitudes that are only half that which was inherent to the cold electron temperature limit. Further nonlinear dynamics developed with the explosive formation of magnetic holes, ending with a saturated state in the form of solitary structures or cnoidal waves. This shows that the incorporation of nonzero temperature results in a weak decrease of their spatial dimensions of the holes and increase of their depth.

*Electronic address: istomin@lpi.ru

†Electronic address: o.a.pokhotelov@sheffield.ac.uk

‡Electronic address: m.balikhin@sheffield.ac.uk

I. INTRODUCTION

The diamagnetic or mirror instability (MI), first theoretically predicted by Vedenov and Sagdeev in Ref. 1, is commonly attributed to the formation of magnetic holes in space plasmas, such as solar wind, planetary magnetosheaths, in the vicinity of comets, the Io wake and the magnetospheric ring current. An increasingly large number of observations have confirmed the existence of the MI in virtually all space plasmas where a proton temperature anisotropy can be generated by some mechanism[2]. The MI generated waves are usually observed in a strongly developed nonlinear state. The study of the nonlinear dynamics of this instability has been a subject of a great deal of research in recent years[3–10]. It should be noted that all previous nonlinear models of the MI were restricted for simplicity by consideration in the cold electron temperature limit which is valid for magnetosheath plasmas. The effect of nonzero electron temperature becomes important whenever $T_e/T_\parallel = O(1)$, where T_e is the electron temperature and $T_\parallel$ is the longitudinal ion temperature. This situation is typical for the solar wind and even exists in magnetospheric plasma where the cold electron model may not be considered as appropriate.

The incorporation of a nonzero electron temperature in the MI linear theory was carried out previously[11,12]. In these papers it was shown that the nonzero electron temperature effect can decrease the growth rate and enhance the instability threshold as well as the angle of wave propagation for the fastest growing mode. Such a modification of the MI is ultimately due to a longitudinal electric field which arises because the electron pressure gradient builds up as the electrons are dragged by the circulating ions from high to low parallel magnetic flux regions.

The incorporation of nonzero electron temperature effects in the nonlinear theory of MI is the main goal of this paper. This will allow us to apply the results of the nonlinear theory not only to the large amplitude mirror waves observed in the magnetosheath but also to the waves observed in other regions of space plasmas (e.g., the ring current).

The paper is organized into the following five sections. In Section II we have derived a set of equations which describe the nonlinear dynamics of MI in the presence of nonzero electron temperature effects. The linearization of these equations in the small amplitude limit is discussed in Section III and the expression for the MI growth rate in the presence of both nonzero electron temperature and finite ion Larmor radius effects is obtained. The temporal evolution of MI and formation of saturated state is analyzed in Section IV. Our discussion and conclusions are found in Section V.

II. BASIC EQUATIONS

Let us consider a low-frequency wave propagating in a plasma immersed in an external magnetic field $\mathbf{B}_0$. A right-handed Cartesian system of coordinates (x, y, z) whose z-axis is directed along the ambient magnetic field is accepted.

We start by considering the perpendicular plasma pressure balance condition[6,8,9]

$$\frac{\delta p_{\perp i} + \delta p_e}{2p^0_{\perp i}} + \frac{1}{\beta_\perp}\left(1 - \frac{3}{4(1+b)^2}\rho_i^2\nabla_\perp^2\right) b + \frac{b^2}{2\beta_\perp}$$

$$= -\frac{1}{\beta_\perp}\left(1 + \frac{\beta_\perp - \beta_\parallel}{2}\right)\nabla_\perp^{-2}\frac{\partial^2}{\partial s^2}b, \tag{1}$$

where $\delta p_{\perp i}$ is the deviation of the ion pressure from its unperturbed state $p^0_{\perp i}$, $b = \delta B_z/B_0$, δB_z the perturbation of the magnetic field along the external magnetic field $\mathbf{B}_0$ direction, $\beta_{\perp(\parallel)} = 2\mu_0 p^0_{\perp(\parallel)i}/B_0^2$ the ion perpendicular (parallel) plasma beta, μ_0 the permeability of free space, $\rho_i = v_{T_\perp}/\omega_{ci}$ the ion Larmor radius in the external magnetic field $\mathbf{B}_0$, $\omega_{ci} = eB_0/m$ the ion cyclotron frequency, e and m the ion charge and mass respectively, $v_{T_\perp} = (2T_\perp/m)^{1/2}$ the perpendicular thermal ion velocity, $T_\perp$ the perpendicular ion temperature, δp_e the perturbation of the electron pressure, and s the distance along the magnetic field $\mathbf{B}$ direction.

The physical meaning of derivation of Eq. (**??**) has already been discussed in our previous paper[9]. In contrast to Ref. 9, in which consideration was limited by the cold electron approximation, Eq. (**??**) includes the variation of the electron plasma pressure δp_e, which for clarity is assumed to be isotropic. Due to their high mobility, the electrons are in an equilibrium state, i.e. their distribution function $f_e \propto \exp(-W/T_e)$, where W is the electron energy $W = m_e v_\parallel^2/2 + \mu_e B - e\Psi$, m_e the electron mass, μ_e the electron magnetic moment, $v_\parallel$ the parallel velocity and Ψ the potential of the electric field. Integrating f_e over $(\mu_e, v_\parallel)$ space one finds that electron density may be described by the Boltzmann law, i.e. $n_e = n_0 \exp(e\Psi/T_e)$ and thus the variation of electron pressure is

$$\delta p_e = n_0 T_e \exp\left(\frac{e\Psi}{T_e}\right), \tag{2}$$

where n_0 is the equilibrium plasma number density and Ψ is related to the field-aligned electric field $E_{\parallel}$ by $E_{\parallel} = -\partial\Psi/\partial s$. The origin of the electric field is the electron pressure gradient created as electrons are dragged from regions of high to low parallel magnetic flux regions.

The ion density and the ion perpendicular pressure are given by the following equations

$$n = B \int f d\mu dv_{\parallel}, \tag{3}$$

and

$$p_{\perp} = B^2 \int \mu f d\mu dv_{\parallel}, \tag{4}$$

where f is the ion velocity distribution function, the subscript i is omitted for clarity, B the magnitude of the total magnetic field, $\mu = mv_{\perp}^2/2B$ the ion magnetic moment, and m the ion mass. For low-frequency oscillations the ion magnetic moment is conserved, i.e. $d\mu/dt = 0$. Furthermore, $v_{\parallel(\perp)}$ is the ion velocity along (perpendicular) the magnetic field lines.

The ion velocity distribution function $f(\mu, v_{\parallel}, \mathbf{r})$, where t is the time and $\mathbf{r}$ the guiding center position vector, in the leading order obeys the drift kinetic equation[9]

$$\frac{\partial f}{\partial t} + v_{\parallel}\frac{\partial f}{\partial s} + \dot{v}_{\parallel}\frac{\partial f}{\partial v_{\parallel}} = 0. \tag{5}$$

The change in the parallel ion velocity resulting from wave compression, $\delta B_{\parallel}$, and an electric field, $E_{\parallel}$, is given by the adiabatic expression[13]

$$\dot{v}_{\parallel} \equiv \frac{dv_{\parallel}}{dt} = -\frac{\mu B_0}{m}\frac{\partial b}{\partial s} - \frac{e}{m}\frac{\partial \Psi}{\partial s}. \tag{6}$$

For clarity we have considered the magnetic and electric field perturbations to be described by harmonic waves in the longitudinal direction, $[b(\mathbf{r},t), \Psi(\mathbf{r},t)] \propto \cos(k_{\parallel}s)$, $|b| < 1$ and $|e\Psi/T_e| < 1$. Introducing the new variables, $2\xi = k_{\parallel}s$ and $\dot{\xi} = k_{\parallel}v_{\parallel}/2$, Eq. (**??**) reduces to the form[9]

$$\frac{\partial f}{\partial t} + \dot{\xi}\frac{\partial f}{\partial \xi} - \frac{\sin 2\xi}{2\tau^2}\frac{\partial f}{\partial \dot{\xi}} = 0, \tag{7}$$

where

$$\tau = \left[\frac{m}{k_\parallel^2 \mu B_0 \, |b| \, (1 + e\Psi / b\mu B_0)} \right]^{1/2}. \qquad (8)$$

The characteristics of Eq. (??) describes the particle motion as a classical nonlinear pendulum whose period T is of the order 2τ. When particles oscillate faster then the wave amplitude grows, i.e. when $T < \gamma^{-1}$ we have an adiabatic regime. Here $\gamma = \partial \ln b / \partial t$ is the instability growth rate. In the opposite case ($\gamma^{-1} < T$) the ions oscillate very slowly in relation to the growth of the wave amplitude. In this limit one can use a linear approximation or polynomial expansion in powers of b.

We note that τ depends not only on the wave amplitude b and the electrostatic potential Ψ but also on the adiabatic invariant μ. Following to Ref. 9 we have introduced the value $\mu = \mu_1 = 4m\gamma^2 / B_0 k_\parallel^2 \, |b| - e\Psi / bB_0$ which separates the regions of adiabatic and nonadiabatic ion motions.

Furthermore, similar to Ref. 9 one finds

$$n = \frac{B}{B_0} n_0 + B \int_0^{\mu_1} d\mu dv_\parallel \delta f, \qquad (9)$$

where δf is the linear perturbation of the ion distribution function given by[14]

$$\delta f = (\mu B_0 b + e\Psi) \left(1 - \frac{\pi \gamma \delta(v_\parallel)}{|k_\parallel|} \right) \frac{\partial F}{m v_\parallel \partial v_\parallel}. \qquad (10)$$

For a bi-Maxwellian distribution the ion number density reduces to

$$n = \frac{B}{B_0} n_0 \left[1 - \frac{T_\perp}{T_\parallel} b \left[\Phi_1(\alpha) + a\Phi_2(\alpha) \right] \left(1 - \frac{\pi^{1/2} \gamma}{|k_\parallel| \, v_{T_\parallel}} \right) \right], \qquad (11)$$

where $v_{T_\parallel}$ is the ion parallel thermal velocity,

$$\Phi_1(\alpha) = 1 - \alpha e^{-\alpha} - e^{-\alpha}, \qquad (12)$$

and

$$\Phi_2(\alpha) = 1 - e^{-\alpha}. \qquad (13)$$

The parameter $\alpha = \mu_1 B_0 / T_\perp$ is given by

$$\alpha = \frac{4\gamma^2 m}{k_\parallel^2 T_\perp |b|} - a, \qquad (14)$$

where a relates the electric potential and magnetic field perturbation, $e\Psi/T_\perp = ab$. Eq. (??) defines the value of α if $4\gamma^2 m/k_\parallel^2 T_\perp |b| - a > 0$. In the case when $4\gamma^2 m/k_\parallel^2 T_\perp |b| - a < 0$ the parameter α is zero. Physically the latter implies that all ions in this case become adiabatic. It was found that in the saturated state $\gamma \to 0$ and $a > 0$.

Taking into account that the electrons are distributed according to the Bolzmann law, $n_e = n_0(1 + abT_\perp/T_e)$, from quasi-neutrality condition $n = n_e$ one finds ($\gamma \ll \left|k_\parallel\right| v_{T_\parallel}$)

$$a = -\frac{T_e}{T_\perp}\left[\frac{(T_\perp/T_\parallel)\Phi_1(\alpha) - 1}{1 + (T_e/T_\parallel)\Phi_2(\alpha)} - \frac{(T_\perp\Phi_1(\alpha) + T_e\Phi_2(\alpha))}{(T_\parallel + T_e\Phi_2(\alpha))^2}\frac{\pi^{1/2}\gamma}{|k_\parallel| v_{T\parallel}}\right]. \tag{15}$$

Similarly, the ion pressure is

$$p_{\perp i} = \frac{B^2}{B_0^2} p_{\perp i}^0 + B^2 \int_0^{\mu_1} \mu d\mu \int_{-\infty}^{\infty} \delta f(\mu, v_\parallel) dv_\parallel, \tag{16}$$

or for a bi-Maxwellian distribution

$$p_{\perp i} = \frac{B^2}{B_0^2} p_{\perp i}^0 - \frac{B_0}{B} b p_{\perp i}^0 \frac{T_\perp}{T_\parallel}\left[\Phi_3(\alpha) + a\Phi_1(\alpha)\right]\left(1 - \frac{\pi^{1/2}\gamma}{\left|k_\parallel\right| v_{T_\parallel}}\right), \tag{17}$$

where

$$\Phi_3(\alpha) = 2 - \alpha^2 e^{-\alpha} - 2\alpha e^{-\alpha} - 2e^{-\alpha}. \tag{18}$$

With the help of Eq. (??) the variation of the total plasma pressure is found to be

$$\frac{\delta p_\perp}{2p_{\perp i}^0} = (1 + \frac{a}{2})b + \frac{b^2}{2} - \frac{b}{2}\frac{T_\perp}{T_\parallel}\left[\Phi_3(\alpha) + a\Phi_1(\alpha)\right]\left(1 - \frac{\pi^{1/2}\gamma}{\left|k_\parallel\right| v_{T_\parallel}}\right). \tag{19}$$

Eqs. (??), (??) and (??) constitute a close set of equations which describe the nonlinear dynamics of the mirror mode perturbations in a plasma with a nonzero electron temperature.

III. THE SMALL WAVE AMPLITUDE LIMIT

The linear regime of the MI corresponds to the large values of α, i.e. $\alpha \gg 1$. Furthermore, in this limit $\Phi_1(\alpha) \to 1$, $\Phi_2(\alpha) \to 1$ and $\Phi_3(\alpha) \to 2$. Therefore, the quasi-neutrality condition (??) reduces to

$$\frac{e\Psi}{T_\perp b} = a = -\frac{T_e}{T_\perp}\left[\frac{T_\perp - T_\parallel}{T_e + T_\parallel} - \frac{T_\parallel(T_\perp + T_e)}{(T_e + T_\parallel)^2}\frac{\pi^{1/2}\gamma}{|k_\parallel| v_{T\parallel}}\right]. \tag{20}$$

which coincides with the corresponding expression of Ref. 12.

It is shown that a takes a nonzero value only when the plasma is anisotropic and the electron temperature is comparable to the parallel ion temperature. Moreover, in the linear MI limit the electric field potential varies in anti-phase with the compressional perturbation of the magnetic field. This may serve as an additional tool for the prime identification of linear mirror perturbations in experimental data.

The linear variation of the total plasma pressure is found from Eq. (**??**)

$$\frac{\delta p_\perp}{2p^0_{\perp i}} = -bA\left(1+\frac{a}{2}\right) + b\frac{T_\perp}{T_\parallel}\left(1+\frac{a}{2}\right)\frac{\pi^{1/2}\gamma}{\left|k_\parallel\right| v_{T_\parallel}}, \tag{21}$$

where $A = T_\perp/T_\parallel - 1$ is the plasma anisotropy.

Substituting Eq. (**??**) into Eq. (**??**) one obtains

$$\frac{\delta p_\perp}{2p^0_{\perp i}} = -b\left(\frac{T_\perp}{T_\parallel} - 1 - \frac{\left(T_\perp/T_\parallel - 1\right)^2 T_e}{2T_\perp(1+T_e/T_\parallel)}\right)$$

$$+b\frac{(1+T_e/T_\parallel)^2 + (1+T_e/T_\perp)^2}{2(1+T_e/T_\parallel)^2}\frac{T_\perp}{T_\parallel}\frac{\pi^{1/2}\gamma}{\left|k_\parallel\right| v_{T_\parallel}}. \tag{22}$$

The physical meaning of the terms on the right-hand side of Eq. (**??**) are as follows: the first two terms in the round brackets represent the plasma pressure anisotropy associated with the ion mirror force whilst the third corresponds to the action of the electrostatic force. In the linear regime these forces act in opposite directions. Finally the last term corresponds to the contribution of resonant ions with small parallel velocities.

Substituting Eq. (**??**) into Eq. (**??**) one obtains the expression for the growth rate of the mirror mode which accounts for both nonzero electron temperature and finite ion Larmor radius effects

$$\gamma = \frac{\left|k_\parallel\right| v_{T_\parallel}}{\pi^{1/2}}\frac{T_\parallel}{T_\perp}\frac{2(1+T_e/T_\parallel)^2}{(1+T_e/T_\parallel)^2 + (1+T_e/T_\perp)^2}\Delta, \tag{23}$$

where

$$\Delta = L - \frac{3}{4\beta_\perp}k_\perp^2\rho_i^2 - \frac{k_\parallel^2}{k_\perp^2\beta_\perp}\left(1+\frac{\beta_\perp-\beta_\parallel}{2}\right), \tag{24}$$

and

$$L = \frac{T_\perp}{T_\parallel} - 1 - \frac{1}{\beta_\perp} - \frac{\left(T_\perp/T_\parallel - 1\right)^2 T_e}{2T_\perp(1 + T_e/T_\parallel)} \equiv A - \beta_\perp^{-1} - E. \tag{25}$$

In the limit when $\rho_i \to 0$ expression (**??**) reduces to that which is obtained in Ref. 12. The incorporation of nonzero electron temperature results in the modification of both the growth rate and instability threshold. In this case the instability appears when $L > 0$, i.e. when $A - \beta_\perp^{-1} > E$. Thus, in the presence of warm electrons the ion anisotropy necessary for the instability onset is greater than that in the cold electron temperature limit. Moreover, for a given ion anisotropy the finite electron temperature effects decrease the instability growth rate. These modifications are ultimately due to the field-aligned electric field which arises in the presence of nonzero electron pressure[11,12]. We note that in the limit of no anisotropy Eq. (**??**) gives damping not growth.

The maximum growth rate is attained when

$$(k_\perp \rho_i)^2_{\max} = \beta_\perp \frac{L}{3}, \tag{26}$$

$$(k_\parallel \rho_i)^2_{\max} = \frac{\beta_\perp^2}{1 + \frac{1}{2}\left(\beta_\perp - \beta_\parallel\right)} \frac{L^2}{12}. \tag{27}$$

The expression for the maximum growth rate now becomes

$$\gamma_{\max} = \frac{\omega_{ci}\beta_\perp L^2}{2\sqrt{3\pi}} \left(\frac{T_\parallel}{T_\perp}\right)^{3/2} \frac{(1 + T_e/T_\parallel)^2}{(1 + T_e/T_\parallel)^2 + (1 + T_e/T_\perp)^2}. \tag{28}$$

From Eqs. (**??**)-(**??**) it follows that $\gamma_{\max}/\left|k_\parallel\right|_{\max} v_{T_\parallel} \propto L \ll 1$. The latter corresponds to the so-called "mirror approximation"[15]. Furthermore, the parameter $(k_\parallel \rho_i)^2_{\max}$ always remains smaller than unity. In order to prove this, the equation (**??**) may be rearranged using the marginal stability condition $L \approx 0$ so that

$$(k_\perp \rho_i)^2_{\max} = \frac{1}{3}\left[1 - \frac{1}{\beta_\perp (A - E)}\right]. \tag{29}$$

Since the onset of the MI corresponds to $\beta_\perp (A - E) > 1$ from Eq. (**??**) it follows that $(k_\perp \rho_i)^2_{\max} < 1/3$. Thus our truncation of the power law expansion of (**??**) in terms of $(k_\perp \rho_i)^2$ is justified.

IV. NONLINEAR EFFECTS

The transition from the linear to the nonlinear regime of MI arises for finite values of α when $\mu_1 \geq T_\perp/B$. Furthermore, Eq. (??) shows that when $\alpha \gg 1$ the parameter a is negative whereas in a strongly nonlinear regime it becomes positive. The reversal of the sign arises when $\Phi_1(\alpha) = T_\parallel/T_\perp$ or when $(\alpha+1)e^{-\alpha} = 1 - T_\parallel/T_\perp$. The actual value of α when the reversal occurs depends on the plasma anisotropy. The smaller the anisotropy, the larger the value of α. For moderate values of the ion anisotropy ($T_\parallel/T_\perp \simeq 1/2$) the reversal arises at $\alpha \simeq 1.65$.

When $T_e \simeq T_\perp$ the value of a is of the order of unity, and transition from linear to nonlinear stage ($\alpha \simeq 1$) occurs for magnetic perturbation amplitude of

$$|b| \simeq |b_L| = 2\gamma_L^2 m_i/k_\parallel^2 T_\perp, \tag{30}$$

which is half that for the case of cold electrons, $T_e << T_\perp$ (cf. Ref. 9).

The latter relation is equivalent to the following differential equation

$$\frac{1}{|b|}\left(\frac{d|b|/dt}{|b|}\right)^2 = \frac{k_\parallel^2 T_\perp}{2m_i}. \tag{31}$$

which results in the explosive temporal growth of $|b|$

$$|b| = |b_L|\left(1 - \frac{t}{t_1}\right)^{-2}, \tag{32}$$

where $t_1 = 2\gamma_L^{-1}$.

Eq. (??) shows that in contrast to the linear regime in which the amplitude grows exponentially, in the nonlinear regime, the amplitude growth becomes explosive. This explosive solution is valid up to the time when the amplitude approaches the stationary value $b = b_s$, when $t = t_1 - \Delta t$, with Δt given by $\Delta t = t_1(b_L/b_s)^{1/2} << t_1$.

It is worth mentioning that in the linear limit the electric potential and compressional magnetic field vary in anti-phase and in the fully developed nonlinear state they are in phase.

Setting $\gamma \to 0$ and $\alpha \to 0$ in Eq. (??) and using the fact that $\Phi_1(\alpha) \to \Phi_2(\alpha) \to \Phi_3(\alpha) \to 0$ one finds

$$\frac{\delta p_\perp}{2p_{\perp i}^0} = \left(1 + \frac{T_e}{2T_\perp}\right) b + \frac{b^2}{2}. \tag{33}$$

Substituting this expression into the perpendicular pressure condition gives

$$\frac{3}{4\beta_\perp}\rho_i^2\nabla_\perp^2 b = (1+\frac{1}{\beta_\perp}+\frac{T_e}{2T_\perp})b + (\frac{5}{2}+\frac{5}{2\beta_\perp}+\frac{T_e}{T_\perp})b^2, \tag{34}$$

which may be rewritten in the dimensionless form as

$$\frac{d^2 b}{dx^2} = \frac{3}{5}\left(\frac{1+\beta_\perp+\beta_e/2}{1+\beta_\perp+2\beta_e/5}\right)b + \frac{3}{2}b^2, \tag{35}$$

where $x = r_\perp/\rho_i[9/20(1+\beta_\perp+2\beta_e/5)]^{1/2}$ and β_e is the electron beta. In the cold electron temperature limit, $\beta_e \to 0$, Eq. (**??**) reduces to the corresponding equation of Ref. 9. The general solution of Eq. (**??**) is the so-called cnoidal wave[9]. In order to understand the basic influence of nonzero electron temperature effects let us consider the particular solution of Eq. (**??**) in the form of solitary wave given by

$$b = -\frac{3}{5}\left(\frac{1+\beta_\perp+\beta_e/2}{1+\beta_\perp+2\beta_e/5}\right)\cosh^{-2}\left[(\frac{3}{20})^{1/2}x\right]. \tag{36}$$

It is easily seen that the nonzero electron temperature shortly decreases the perpendicular size of the magnetic structure. Furthermore, this effect also slightly increases the depth of the magnetic hole. With the growth of electron temperature the dimensionless depth of the magnetic hole varies from $-3/5$ when $\beta_e << \beta_\perp$ to $-3/4$ when $\beta_e >> \beta_\perp$, i.e. from 60 to 75 percent of the ambient magnetic field.

V. SUMMARY

In this paper, we have presented a study of nonlinear dynamics of the MI in the presence of nonzero electron temperature effects using direct integration of the drift-kinetic equation for the ions and a Boltzmann distribution for the electrons. A dynamic model was then developed which accounts for the field-aligned electric field existing in a system with nonzero electron pressure. The main characteristic of the present model is the role of the electrostatic force both in the linear and nonlinear regimes. It was found that in the linear approximation the electrostatic potential is in anti-phase with the variation of the compressional magnetic field perturbation. This contrasts with the saturated state in which they are in phase. It has been found that in the saturated state the nonzero electron temperature effects slightly decreases the size and increases the depth of the magnetic structure. These changes become noticeable only when the electron temperature becomes comparable with the parallel ion temperature.

Recently further progress in the study of nonlinear MI was carried out[16]. The authors of this paper found two different solutions in the form of magnetic humps and holes, resulting from the wave-wave and wave-particle coupling, respectively born on the same physical conditions. However, nonzero electron temperature effects were not included in this study.

The model developed in our paper still remains oversimplified. For example, it has been restricted to the case of isotropic electrons and finite but relatively small amplitudes of the solitons when $|b| < 1$ and the Korteweg-de-Vries expansion provides a useful guide for the construction of the nonlinear equations. The case when $\delta B \sim B$ was considered in Ref. 16. Furthermore, the effect of bistability of mirror modes revealed in recent observations and discussed in Refs. 7 and 16 was also outside the scope of this current study. However, our analysis has provided a deeper insight into the physics of the nonlinear dynamics of mirror modes in high-β space plasmas.

Acknowledgments

The work was supported by the STFC, by the Russian Fund for Basic Research grants No 08-05-00617, by Programs 7 and 16 of the Presidium of the Russian Academy of Sciences, by ISSI and by the ISTC project No 3520.

[1]A. A. Vedenov and R. Z. Sagdeev, in *Plasma Physics and Problem of Controlled Thermonuclear Reactions*, edited by M. A. Leontovich (Pergamon, NY, 1958), Vol. III, p. 332.

[2]J. Soucek, E. Lucek, and I. Dandouras, J. Geophys. Res., **113**, A04203, doi:10.1029/2007JA012649 (2008).

[3]F. G. E. Pantellini, D. Burgers and S. J. Schwartz, Adv. Space Res., **15**, 341 (1995).

[4]M. G. Kivelson and D. J. Southwood, J. Geophys. Res., **101**, 17365 (1996).

[5]F. G. E. Pantellini, J. Geophys. Res., **103**, 4789 (1998).

[6]E. A. Kuznetsov, T. Passot, and P. L. Sulem, Phys. Rev. Lett., **98**, 235003 (2007).

[7]F. Califano, P. Helliger, E. Kuznetsov, T. Passot, P. L. Sulem, and P. M. Trávniček, J. Geophys. Res., **113**, A08219, doi:10.1029/2007JA012898 (2008).

[8]O. A. Pokhotelov, R. Z. Sagdeev, M. A. Balikhin, O. G. Onishchenko and V. N. Fedun, J. Geophys. Res., **113**, A04225, doi:10.1029/2007JA012642 (2008).

[9]Ya. N. Istomin, O. A. Pokhotelov, and M.A. Balikhin, Phys. Plasmas, **16**, 062905, doi:10.1063/1.3153553 (2009).

[10]P. Hellinger, E. A. Kuznetsov, T. Passot, P. L. Sulem, and P. M. Trávniček, Geophys. Res. Lett., 36, L06103, doi:10.1029/2008GL036805 (2009).

[11]F. G. E. Pantellini and S. J. Schwartz, J. Geophys. Res., **100**, 3539, 1995.

[12]O. A. Pokhotelov, M. A. Balikhin, H. St-C.K. Alleyne, and O. G. Onishchenko, J. Geophys. Res., 105, 2393 (2000).

[13]T. G. Northrop, *The Adiabatic Motion of Charged Particles* (Wiley Interscience, New York, 1963).

[14]O. A. Pokhotelov, R. A. Treumann, R. Z. Sagdeev, M. A. Balikhin, O. G. Onishchenko, and V. P. Pavlenko, J. Geophys. Res., **107**, 1312, doi:10.1029/2001JA009125 (2002).

[15]Southwood, D. J. and M. G. Kivelson, J. Geophys. Res., **98**, 9181 (1993).

[16]D. Jovanović, and P. K. Shukla, Phys. Plasmas, **16**, 082901, doi: 10.1063/1.3183591 (2009).