Probing Sgr A* with theory of large masses without events horizon

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Abstract

In the present paper some consequences of the assumption that in the center of the Galaxy there is a supermassive compact object without the events horizon are considered. The possibility of existence of such object has been argued earlier. It is shown, that accretion of a surrounding gas onto the object can cause nuclear burning in a superficial layer which owing to comptonization in a hotter layer, laying above, can manifest itself in observable IR and X spectra. The contribution of an intrinsic magnetic moment of the object in the observable synchrotron radiation is considered, using transfer equations, taking into account influence of gravitation on the energy and movement of photons.
1 Introduction

An analysis of stars motion in the dynamic center of the Galaxy give strong evidence for the existence of a compact object with mass about $3 \cdot 10^6 M_\odot$ or more that is associated with Sgr A* [1,2,3,4]. There are three kinds of an explanation of observable peculiarities of the object radiation:

1 - The gas accretion onto the central object – a supermassive black hole (BH) [5,6].

2 - An ejection of the magnetized plasma from the vicinity of the Schwarzschild radius of the BH [7,8].

3 - Explanations based on hypotheses about another nature of the central object (a cluster of dark objects [9], a fermion ball hypothesis [10], boson stars [11]).

In the present paper we consider some consequence of the assumption that radiation of Sgr A* is caused by existence of a supermassive compact object without events horizon in the Galaxy Center. Such steady configurations of the degenerated Fermi-gas with masses $10^2 \div 10^5 M_\odot$ and with the radii $R$ less than the Schwarzschild radius $r_g$ are one of the consequence of the metric-field equations of gravitation [12], [13], [14]. In the theory gravitational field of an attractive mass manifests itself as a field in Minkowski space-time for a remote observer in an inertial frame of reference, and as space-time curvature for the observer in a co-moving (with the free falling particles) frame of reference. Physical consequences from the gravitation equations under consideration are very close to the ones in general relativity at the distances from the central mass much more than $r_g$. However they are completely different at the distances nearby $r_g$ or less than that. The spherically-symmetric solution of the gravitation equations have no the event horizon and physical singularity in the center [16].

Since the gravitation equations was tested by the binary pulsar PSR 1913+16 [15] and stability of the supermassive configurations was studied sufficiently rigorously [13], it is meaningful to investigate the possibility of the existence of such objects at the Galaxy Center as an alternative to the supermassive black hole hypothesis.

The gravitational force of a point mass $M$ affecting a free falling particle of mass $m$ is given by [16]

$$F = -m \left[ \frac{c^2 B_1 + (B_2 - 2B_3) r^2}{r^2} \right],$$  \hspace{1cm} (1)

where

$$B_1 = C'/2A, \quad B_2 = A'/2A, \quad B_3 = C'/2C$$  \hspace{1cm} (2) \hspace{1cm} \text{and}

$$A = f^2/C, \quad C = 1 - r_g/f, \quad f = (r_g^3 + r^3)^{1/3}.$$  \hspace{1cm} (3)

In this equation $r$ is the radial distance from the center, $r_g = 2GM/c^2$, $M$ is the mass of the object, $G$ is the gravitational constant, $c$ is speed of light, the prime denotes the derivative with respect to $r$. The force affecting the test particle in rest is

$$F = -\frac{GmM}{r^2} \left[ 1 - \frac{r_g}{(r^3 + r_g^3)^{1/3}} \right]$$  \hspace{1cm} (4)

Fig. 1 shows the force $F$ (in arbitrary units) affecting the test particle at rest (curve 1) and the free falling particle (curve 2) as the function of the distance $r/r_g$ from the center.

It follows from the above plot that the gravitational force affecting free falling particles decreases when $r$ approach to $r_g$ and changes its sign at $r = 1.5 r_g$. Although we never observed the motion of the particle at distances close to $r_g$, we can test this conclusion for very distant objects in our Universe because for its observed mass $M_u$ the value of $2G M_u/c^2$ is close to the observed radius $r_u$ of the Universe. At such distances the gravitational force affecting the particles change the sign. And the repulsion force in a simple model of the expanded selfgraviting dust ball gives a simple and clear explanation of the acceleration of the Universe expansion [17], [18].

It seems in the first sight that the accretion onto the object give rise to a too large energy release at the surface that contradicts the low bolometric luminosity ($\sim 10^{36}$ erg s$^{-1}$) of Sgr A*. However, it must be taken into account that in the gravitation theory under consideration the velocity of free falling test particles decrease inside the Schwarzschild radius [16]. If we assume that the the radius $R$ of the object with mass of $2.6 \cdot M_\odot$ in the Galactic Center is equal to $0.04 r_g$ (which follows from the solution of the equation of the hydrostatic equilibrium [12], [13], [14]), then the value of the velocity $v$ of free falling particles at the surface is $4 \cdot 10^8$ cm s$^{-2}$. Therefore, even at the accretion rate $\dot{M} = 10^{-6} M_\odot$ yr$^{-1}$ the amount of the released energy $\dot{M} v^2/2$ is only $\sim 10^{36}$ erg/s.
2 Atmosphere

It is believed that the rate of gas accretion onto the supermassive object in the Galaxy center due to the star wind from the surrounding young stars is of the order of $10^{-7} M_{\odot} \text{ yr}^{-1}$ [12]. Therefore, if the object has a surface, it must have also some, mainly hydrogen, atmosphere. For $10 M_{\odot} \text{yr}$ (an estimated lifetime of the surrounding stars) the mass $M_{\text{atm}}$ of the gas envelope reaches $M_{\odot}$. The height of the homogeneous atmosphere $h_{\text{atm}} = kT/2 m_p g$, where $k$ is the Boltzmann constant, $T$ is the absolute temperature, $m_p$ is the proton mass, $g = F/m = 7.8 \cdot 10^6 \text{ cm s}^{-2}$. At the temperature $T = 10^7 K$ the high of the atmosphere $h_{\text{atm}} = 10^8 \text{ cm} \approx 0.01 R$. The density of the atmosphere $\rho = M_{\text{atm}}/4\pi R^2 h \approx 10^2 \div 10^3 g \text{ cm}^{-3}$. Under the condition a hydrogen burning must begin in our time. Of course, the accretion rate in the past it could has been many order greater than the mention above if the stars was born in a molecular cloud at the same area where they are in our time. (See the discussion in [11,13].) In this case the burning began more than $10 M_{\odot} \text{yr}$ ago and is more intensive.

A relationship between the temperature and the density can be found from the thermodynamics equilibrium equations

$$\frac{1}{4\pi r^2 \rho} \frac{dL_r}{dr} = \epsilon + \lambda \frac{M c^2}{M_{\text{atm}}} - T \frac{dS}{dt},$$

where

$$L_r = 4\pi r^2 \frac{ac}{3\kappa \rho} \frac{dT}{dr} r^2,$$

$M$ is the rate of the mass accretion, $\epsilon$ is the rate of the generation of the nuclear energy per the unit of mass, $\lambda$ is the portion of the thermalized accretion energy, $\kappa$ is the Krammers absorbtion factor, $S$ is entropy, $a$ is the radiation constant. In the stationary case for the homogeneous atmosphere we have

$$\frac{ac T^4}{3 \kappa \rho^2 h_{\text{atm}}^2} = \epsilon + \lambda \frac{M c^2}{M_{\text{atm}}},$$

Fig. 2 shows the relationship between $T$ and $\rho$ for several values of the parameter $\lambda$. It follows from the figure that the hydrogen burning can be occurs at the density $10^2 \div 10^3 g \text{ cm}^{-3}$.

The luminosity from the burning layer can be found by solution of the system of the differential equations

$$\frac{dm_r}{dr} = 4\pi r^2, \quad \frac{dP}{dr} = \rho g, \quad \frac{dL_r}{dr} = 4\pi r^2 \rho c, \quad \frac{dT}{dr} = -\frac{3\kappa \rho L_r}{16 \pi a c r^2 T^3},$$

where $P = \rho kT/2 m_p$ is the pressure of the nogenerated gas, $m_r$ is the mass of the spherical layer of the radius $r$. The boundary data at the surface are: $m_r(R) = 0$, $\rho(R) = \rho_0$, $T(R) = T_0$, $L_r(R) = 0$. We setting $T_0 = 10^7 K$ and find that the luminosity $L = 2.4 \cdot 10^{31} \text{ erg s}^{-1}$ at $\rho_0 = 10^2 g \text{ cm}^{-3}$ and $L = 2.4 \cdot 10^{33} \text{ erg s}^{-1}$ at $\rho_0 = 10^2 g \text{ cm}^{-3}$.

The effective temperature $T_{eff}$ is $2.4 \cdot 10^3 K$ and $7.8 \cdot 10^3 K$, respectively. These magnitudes may be much more if in the past the accretion was more intensive. However, if this is not true, and if the accretion rate is much less than $10^{-7} M_{\odot}$, the hydrogen burning presently may be missing.

3 Peculiarity of the accretion

The main peculiarity of a spherical supersonic accretion onto the supermassive object without events horizon is the existence of the second sonic point – in a vicinity of the object, which does not bound with hydrodynamical effects. The physical reason is that as the sound velocity $v_s$ in the infalling flow grows together with the temperature, the gas velocity $v$ begin to decrease fast before $r = r_g$ [16]. As a result, the equality $v = v_s$ take place at some distance $r_s$ from the center.

In accordance with [16] the maximal radial velocity of a free falling particle does not exceed $0.4 c$. Therefore, the Lorentz-factor is nearly 1 and the accretion equations are given by

$$4\pi r^2 v = \dot{M}, \quad vv + \dot{n}/n = g_{fr}$$

$$\left( \frac{\varepsilon}{n} \right)' - P \frac{n'}{n} = -\frac{\Lambda}{v n}.$$  

In these eqs. $n$ is the density of the particles number, $\Lambda$ is the cooling rate (per volume unit) due to the bremsstrahlung and comptonization, $g_{fr} = F/m$ is defined by eq. [1]. We, therefore, neglect the radiation pressure and do not take into account the subequispartion magnetic field which may exist in the accretion flow [5].

The energy density is assumed to be equal to [5]

$$\varepsilon = m_p c^2 n + \alpha n kT,$$
where
\[ \alpha = 3 + x \left( \frac{3K_3(x) + K_1(x)}{4K_2(x)} - 1 \right) \quad (11) \]
\[ x = m_e e^2 / kT, \quad y = m_p e^2 / kT \] and \( K_j \) (j = 2, 3) is the modified Bessel function.

We assume for definiteness that \( \dot{M} = 10^{-7} M_\odot \ yr^{-1} \) and at the distance from the center \( r_0 = 10^{17} \ cm \) the velocity \( v(r_0) = 10^8 \ cm \ s^{-1} \) and the temperature \( T(r_0) = 10^5 \ K \).

The profile of the temperature and velocity can be found by the solution of the eqs. \[22\] The sound velocity \( v_s = v_s(r) \) is
\[ v_s = \left( \frac{\Gamma P}{\rho + \rho P} \right)^{1/2}, \quad (12) \]
where \( \rho \) is the density, the adiabatic index is
\[ \Gamma = \frac{n}{P} \left( \frac{\partial P}{\partial n} \right) S_p \quad (13) \]
and \( S_p \) is entropy per particle. The values of \( P, \rho \) and \( \Gamma \) for a perfect Boltzmann gas as the function of \( T \) is given by the Service fitting formulas \[20\]. Under these conditions the gas temperature \( T \) at \( r = r_s \) is of the order of \( 10^{10} \ K \). The postshock region lies from \( r_s \) up to the dense atmosphere ( \( n \approx 10^{15} \div 10^{17} \) ) at the distance \( z \approx R \) from the object surface where infalling gas velocity slumps. In this area releases more part of the accretion energy. (We neglect here a difference between the height of proton and electron stopping \[29\]. The energy release per 1 g due to protons stopping is given by \[22\]
\[ W = -f(x_e) \frac{4\pi n e^4 \log \Delta}{m_e c m_p}. \quad (14) \]

where
\[ \Delta = \frac{3(kT)^{3/2}}{4\sqrt{\pi} n e^3} \] : \[ x_e = \left( \frac{m_e v^2}{kT} \right)^{1/2} \quad (15) \]
and
\[ f(x_e) = \frac{2x_e(2x_e^2 - 3 - m_e/m_p)}{1 + 2x_e(2x_e^2 - 3 - m_e/m_p)}^{3/2}. \quad (16) \]

We assume that the gas density profile in this area is the sum of the atmosphere density resulting from eqs. \[3\] and the solution \( n(r) \) of the eqs. \[4\]. Let \( U \) be the radiation density and \( k_2 = 0.4 + 6.4 \cdot 10^{22} m_p n T^{-7/2} e^2 cm^{-3} \) is the flux mean opacity. At the given density the profile ("jump") of the luminosity and temperature can be found approximately from the equations
\[ dL_r / dr = 4\pi r^2 \rho W, \quad \frac{1}{3 \rho} \frac{dU}{dr} = k_2 \frac{L_r}{4\pi r^2 c}, \quad (17) \]

and from the simple energy balance equation [Zeldovich et al. 1969] between \( W \), bremsstrahlung and comptonization processes. Fig. 3 shows a typical temperature profile close the distance \( z = R \) from the object surface. The following end conditions were used: the luminosity at the distance from the surface 0.0785R = 10^{32} \ erg \ s^{-1}, \ the temperature is 10^7 \ K, the density of radiation is \( a T_{eff}^4 \), where \( T_{eff} = 10^4 \ K \) and \( a = 7.569 \cdot 10^{15} \).

4 Comptonization

The emerging intensity \( I \) of the low atmosphere radiation after passage through a hot spherical homogeneous lay of gas is a convolution of the incoming Plankian intensity \( I_0 \) and the frequency redistribution function \( \Phi(s) \) of \( s = \log (\nu/\nu_0) \), where \( \nu_0 \) and \( \nu \) are the frequencies of the incoming and emerging radiation, respectively:
\[ I(x) = \int_{-\infty}^{\infty} \Phi(s) I_0(s) ds. \quad (18) \]

At the temperature of the order of \( 10^{10} \ K \) the dimensionless parameter \( \vartheta = kT/m_e e^2 \sim 1 \) or more that that. Under the condition the function \( \Phi(s) \) can be calculated as \[21\]
\[ \Phi(s) = \sum_{k=0}^{m} \frac{e^{-\tau k} \tau^k}{k!} P_k(s), \quad (19) \]

where
\[ P_k(s) = \int_{-1}^{1} \varphi(\beta) P(s, \beta) d\beta, \quad (20) \]
\[ \beta_{min} = \frac{e^{\vartheta} - 1}{e^{\vartheta} + 1} \quad (21) \]
and
\[ P(s, \beta) = \frac{3}{16\gamma^2 \beta} \int_{\mu_1}^{\mu_2} (1 - \beta \mu)^{-3} \left( 1 + \beta \mu' \right) \left( 1 + \mu \mu' \right) \left( 1 - \mu^2 \right) \left( 1 - \mu'^2 \right) d\mu, \quad (22) \]
Figure 4: The emerging spectrum of the nuclear burning after comptonization by the hot layer at the optical thickness $\tau = 0.01, 0.1, 1, 5$

and $\gamma = (1 - \beta^2)^{-1/2}$. Taking into account the scattering up to 4th order ($m = 4$) we find an approximate emerging spectrum of the nuclear burning with $T_{eff} = 10^4 K$ after comptonization for several values of the optical thickness $\tau$ of a homogeneous hot layer. The results are plotted in fig. 4.

It follows from the figure for small magnitudes of $\tau$ that the spectrum can contribute to the observed Sgr A* radiation \cite{23}, \cite{24}, \cite{25} in IR and X regions.

It should be noted that the timescale $\Delta t$ of a variable process are related to the size $\Delta r$ of its region as $\Delta r \leq c \Delta t/(1 + z_g)$, where $1 + z_g = 1/\sqrt{C(r)}$ is gravitational redshift. (The function $C(r)$ is defined by eqs. 3).

Therefore, the variations in the radiation intensity $\sim 600 s$ happen at the distance $z \leq 1.7R$ from the surface. For this reason IR and X flares \cite{23}, \cite{24}, \cite{25} can be interpreted as processes near the surface of the objects.

5 Transfer equation and synchrotron radiation of Sgr A*

Some authors \cite{26}, \cite{27} have found evidence for existence of an intrinsic magnetic field of Black Holes candidates that is incompatible with the events horizon. Here we calculate the contribution in the synchrotron radiation spectrum from the intrinsic magnetic field of the form

$$B = B_0 \left(\frac{R}{r}\right)^3, \quad (26)$$

where $B_0 = 2 \cdot 10^5 Gs$. For study of the radiation from the surface vicinity of the supermassive object the influence of gravitation on the frequency of the photon and its motion must be taken into account. A relativistic transfer equation can be used for this purpose. It is a relativistic Boltzmann equations for photon gas \cite{Lindquist 1966}, \cite{Smidt-Burgk}, \cite{Zane et al. 1996}. We assume that in the spherically symmetric field the distribution function $F$ depends on the radial distance from the center $r$, the frequency $\nu$ and the photon direction which can be defined by the cosine of the horizontal angle $\mu$. (The spherical coordinate system is used). The relativistic Boltzmann equation in Minkowski space - time is of the form

$$\frac{dF}{dn} = St(F), \quad (27)$$

where $dF/dn$ is the derivative along the 4-trajectory $x^\alpha = x^\alpha(s)$ in space-time with the metric differential form $dn$ and the right-hand member is the collisions integral. By using for photons $x^0 = ct$ as a parameter along 4-trajectory we arrive at the transfer equation in the form

$$\left( \frac{\partial}{\partial dx^0} + n\nabla + \frac{d\mu}{dx^0} \frac{\partial}{\partial \mu} + \frac{d\nu}{dx^0} \frac{\partial}{\partial \nu} \right) F = St(F), \quad (28)$$

where the magnitudes $d\mu/dx^0$ and $d\nu/dx^0$ must be found from our equations of the photon motion in the gravitation field \cite{12}. The collisions integral is given by

$$\chi (S/\beta - F), \quad (29)$$

where $\chi$ is the absorption coefficient, $S = \eta/\chi$ is the source function, $\eta$ is the emissivity and $\beta = v^2/c^2$. In this paper we do not take into account light diffusion. The intensity $I$ of the radiation is related to $F$ as $I = \beta F$.

To solve the transfer equation we use the characteristic method \cite{Smidt-Burgk}, \cite{Zane et al. 1996}. Since photon’s trajectories are characteristics of the partial differential equation \cite{28}, the equations are reduced to ordinary differential equations along these trajectories. In our case this equation is

$$\frac{dF}{dr} = \frac{c}{v_{ph}} \left( \frac{S}{\beta} - F \right), \quad (30)$$

where $v_{ph}$ is the radial photon velocity.

According to \cite{16} in the spherically - symmetric field

$$v_{ph} = c \sqrt{\frac{C}{A} \left( 1 - \frac{b^2}{f^2} \right)}, \quad (31)$$

where $b$ is the impact parameter of photon.

For numerical estimates we assume that a relativistic Maxwellian electron distribution take place and set emission coefficient \cite{28}

$$j_\nu = \frac{2^{1/6} \pi^{3/2} e^2 n_e \nu}{3\sqrt{2} c K_2 (1/\Theta) 1/6} \nu^{3/2} \exp[-(9\nu / 2)^{1/3}], \quad (32)$$
Figure 5: The function $b(r)$

where $\Upsilon = \nu / \nu_c \Theta^2$ and $\Theta = kT/mc^2$, $n_e$ is the electron density number, $K_2$ is the modified Bessel function of the second kind.

Let $\mathcal{F}_\nu(b)$ be the solution of the differential equation \ref{eq:30} for a given $b$ at $r \to \infty$. Then, for a distant observer the luminosity at the frequency $\nu$ is given by

$$L_\nu = 8\pi^2 \int_0^\infty \beta \mathcal{F}_\nu(b) db. \tag{33}$$

For a correct solution of eq.\ref{eq:30} it is essential that there are three types of photons trajectories in the spherically-symmetric gravitation field in view the used gravitation equations. It can be seen from fig. 5. It shows the geometrical locus where the radial photon velocities are equal to zero which is given by the equation

$$b = \frac{f}{\sqrt{C}}\tag{34}$$

The minimal value of $b$ is $b_{cr} = 3\sqrt{3}r_g/2$. It occurs at the distance from the center $r_{cr} = \sqrt{19}r_g/2$. The photons whose impact parameter $b < b_{cr}$ can freely move from the object surface to infinity. The photons with $b > b_{cr}$ cannot go away from the surface to infinity. At last, photons with $b > b_{cr}$ can move to infinity only if their trajectories begin at distances $r > r_{cr}$. For a given $b$ the corresponding magnitude of $r$ can be found from eq.\ref{eq:34}

The differential equation for photon trajectories with $b < b_{cr}$ were integrated at the edge condition $\mathcal{F}(R) = 0$. The distribution function $\mathcal{F}$ at the points of the curve in fig. 5 at $r > r_{cr}$ was found by solution of differential equation \ref{eq:30} by used the end condition $\mathcal{F} = 0$ at infinity. Similarly to \cite{Zane et al. 1996} these magnitudes were used as end conditions for the solution of the eq. \ref{eq:30} to find the emergent radiation.

Fig. 6 shows the spectrum of the synchrotron radiation caused by the proper magnetic field of the object.

- for the sum of the above magnetic fields (thick curve).
- It follows from the figure that existence of the intrinsic magnetic field may give a good fitting to the observation data.

6 Conclusion

The assignment of nature of compact objects in the galactic centers is one of basic problems of fundamental physics and astrophysics. The results received above, certainly, yet do not allow to draw the certain conclusions. However they show that the opportunity investigated here does not contradicts the observant data, and, therefore, demands the further study.

References

[1] Genzel, R.; Schödel, R.; Ott, T. et al. : 2003, ApJ 594, 812.
[2] Schödel,R., Ott, T., Genzel, R.: 2003, astro-ph/0306214.
[3] Ghes, A.M., Salim, S., Hornstein, S.D., Tanner, A., Morris, M., Becklin, E.E., Duchene G. astro-ph/0306130.
[4] Ghes, A.M., Duchêne, G., Matthews, K. et al.: 2003, ApJ 586,127.
[5] Coker, R.F., Melia, F.: 2000, ApJ 534, 723.
[6] Narayan,R., Yi,I.: 1995 ApJ 444, 231.
[7] Falcke, H., Markoff, S.:2000, A&A 362,113.
[8] Melia, F., Falcke, H.:2001, ARA&A 39,309.
[9] Maoz,E.:1998, ApJ 494, 181.
[10] Viollier, R..:2003, AN 324,50.
[11] Torres, D.F., Capozziello, S., Lambiase, S.: 2000, PhysRev D 62 104012.
[12] Verozub, L.V.: 1996, AN 317, 107.
[13] Verozub, L.V., Kochetov, A.Ye.: 2001, AN 322, 143.
[14] Verozub, L.V.: 2001 In: "Gravitation, Cosmology and Relativistic Astrophysics", Kharkov Univ., 44.
[15] Verozub, L.V., Kochetov, A.Ye.: 2000, Grav and Cosm. 6, 246.
[16] Verozub, L.V.: 1991, Phys.Lett.A 156, 404.
[17] Verozub, L.V.: 2002 AIP Conference Proc. 624, 371.
[18] Verozub, L.V.: 2002 In: "New Trends in Theoretical and Observational Cosmology", (Ed. by K. Sato & T. Shiromizu), Univ. Acad. Press, Tokyo 2002, 371.
[19] Coker, R.F.: 2001, A&A 375, 18.
[20] Service, A.T.: 1986, ApJ 307, 60.
[21] Birkinshaw, M.: 1999, Physics Reports 310, 97.
[22] Alme, M.L., Wilson, J.R.: 1973, ApJ 186, 1015.
[23] Baganoff, F.K., Maeda, Y., Morris, M. et al.: 2003, ApJ 591, 891.
[24] Porquet, D., Predehl, P., Aschenbach, B. et al.: 2003.
[25] Genzel, R., Schödel, R., Ott, T. et al.: 2003, Nature 425, 934.
[26] Mitra, A.: 2002, Bull. Astron. Soc. India 30, 173.
[27] Robertsob, S.L., Leiter, H.: 2001, ApJ 515, 365.
[28] Wardziński, G., Zdziarski, A.: 1999, MNRAS 314, 183.
[29] Bieldsten, L., Salpeter, E., Wasserman, I.: 1992, ApJ 384, 143.
[Lindquist 1966] R.W. Lindquist, R.W.: 1966, Ann. Phys. 37, 487.
[Smidt-Burgk] Schmid-Burgk, J.: 1978, APSS 56, 191.
[Zane et al. 1996] S. Zane: 1996, astro-phys/9602034.
[Zeldovich et al. 1969] Zeldovich, Ya. B., Shakura, N. I.: Astron. J. (In Russian), 46, 2, 225, 1969.