Solid State Communications, Vol. 55, No. 2, pp. 187–191, 1985.
Printed in Great Britain.

0038–1098/85 $3.00 + .00
Pergamon Press Ltd.

ON THE POSSIBLE SUPERFLUIDITY OF BIPOLARONS ON THE JUNCTION SURFACE

S. Brazovskii, N. Kirova and V. Yakovenko

L.D. Landau Institute for Theoretical Physics. Moscow, USSR

(*Received* 1 *February* 1985 *by V.M. Agranovich*)

There is a number of materials where bipolarons (BP) are probably the main extrinsic carriers. We discuss a dense system of BP produced by the equilibrium charge injection at the surface of Schottki or MIS junctions. Possible super-conductivity of this system due to the Bose-condensation of BP is studied.

1. INTRODUCTION

RECENTLY THERE HAVE BEEN obtained a number of experimental and theoretical results suggesting the existence of bipolarons (BPs) in some materials. The most complete data have been obtained for the compound Ti_4O_7 where BPs exist in the form if diatomic centres [1]. The compound $WO_{3-\delta}$ also has been considered from this point of view [2].

The data on the absence of spin paramagnetism and some optical data for conducting dopped polymers (polyphenylene [3], polypyrrole [4] and polythophene [5]) also suggest the presence of BPs in these materials. The existence of BPs in quasi-one-dimensional materials has theoretically been predicted for Peierls dielectrics (e.g. *trans*-polyacetylene) with weak [6] or strong [7] interchain couplings and for systems with the so called "combined Peierls state" (e.g. for *cis*-polyacetylene [8]).

At present it has become typical that BPs are observed om quasi-one-dimensional semiconductors which is quite natural since these materials are characterised by deep selftrapping and by high values of the dielectric constants. The general theoretical conclusions for the materials like polypyrrole and polyphenilene have been confirmed by the computer calculations [9] and by the exact solutions of the structural models [10].

The most interesting and convincing manifestation of BPs would be their superfluidity, i.e. superconductivity. This phenomenon may be realised in principle because the Bose-condensation in a system of charged bosons is possible [11]. The detailed theory of the Bose-condensation and crystallisation of BPs in the model for Ti_4O_7 but without the Coulomb interactions has been developed in [12]. Nevertheless manifestation of BP's superconductivity in the existing homogeneous materials does not seem to be plausible. The problem is that in all known cases there are no intrinsic mechanisms for the charge transfer, which would fix a finite concentration of BP's in a volume. In principle the concentration of BPs can be maintained by doping. (In the case of $V_xTi_{4-x}O_7$ where at low temperatures BPs occupy all available sites the doping by Vanadium creates BPs' vacancies). But here the random doping can suppress the Bose-condensation by disorder effects which are not yet theoretically studied. (A possible exception can be found in polypyrrole and polythiophene where the chains form well defined planes and the doping is actually intercalation). Some materials, being theoretically promising candidates for the BPs' existence like *cis*-polyacetylene and polydiacetylenes, can not be dopped at all. For this reason we suggest considering an inhomogeneous concentration of BP's at the surface of the sample near the junction with the metal or in MOS structure. To clarify the general possibility of the BPs' Bose-condensation at the junction we have to study a number of obstacles.

1. The Bose-condensation is usually considered for systems with a fixed number of particles which is not true in our case.

2. The Bose-condensation of the two-dimensional ideal gas is not possible in any sense.

3. The condensation of the nonideal gas can be possible in the sense of Beresinskii, Kösterlitz, Thouless (BKT) but the influence of the Coulomb forces is unclear.

The problem of the Coulomb force influence has already been dealt within the theory of superconducting films.

2. THE GROUND STATE

Let Δ be one-electron energy prior to lattice polarisation i.e. $\Delta = E_{opt}/2$, E_{opt} being the optical gap. The energy of a stationary one-electron state (a polaron) is $E_1 = \Delta - \epsilon_1$, $\epsilon_1 > 0$. The energy of a stationary two-electron bound state (a BP) is $E_2 = 2\Delta - \epsilon_2$, $\epsilon_2 > 0$. We are interested in the case when BPs are the main extrinsic carriers (EC), i.e. when they provide the lowest energy per an electron: $\epsilon_2/2 > \epsilon_1$ (e.g. in a model of shallow selftrapping without the Coulomb force taken into

account we have $\epsilon_2 = 4\epsilon_1$). At the same time we can anticipate that $E_2 > E_1$, i.e. $\Delta > \epsilon_2 - \epsilon_1$. These two conditions imply that the main thermoactivated (intrinsic) carriers (IC) are polarons whereas the main EC are BPs.

We shall consider only the case of low temperatures T when we can neglect all activated carriers. Then the system is described by the Lagrangian:

$$\mathcal{L} = \int_0^\infty dz \int d^2r \left\{ \psi^* \left[\frac{\partial}{\partial \tau} + \frac{1}{2M_\parallel} \frac{\partial^2}{\partial r^2} + \frac{1}{2M_\perp} \frac{\partial^2}{\partial z^2} + E_2 \right] \psi \right.$$
$$+ \phi(\mathbf{R})[2en(\mathbf{R}) + 2e'n(\mathbf{R}')]$$
$$\left. - \frac{1}{8\pi} \left[\epsilon_\perp \left(\frac{\partial \phi}{\partial z} \right)^2 + \epsilon_\parallel \left(\frac{\partial \phi}{\partial \mathbf{r}} \right)^2 \right] \right\}, \qquad (1)$$

where $0 < \tau < \beta = T^{-1}$, $\mathbf{R} = (\mathbf{r}, z)$, $\mathbf{R}' = (\mathbf{r}, -z)$, $\mathbf{r} = (x, y)$ and z are coordinates in the plane of the junction ($z = 0$) and normal to it, respectively, $\epsilon_\parallel$, $\epsilon_\perp$ are the dielectric constants for the substance at $z > 0$, $M_\parallel$, $M_\perp$ are effective masses of BPs, e is the electonic charge, $2e'$ is the image charge, ϕ is the Coulomb potential and ψ is the Bose-field operator.

Two cases should be considered. For contact with a dielectric we fix the values of the electric displacement $D \pm \infty$ at $z \to \pm \infty$

$$\frac{1}{4\pi} \{D(z = -\infty) - D(z = +\infty)\} = \int_0^\infty dz\, |\psi|^2 = \sigma,$$

where σ is the BPs' surface density.

For a contact with a metal we fix the BPs' chemical potential $\mu_2 = 2\mu_0$, where μ_0 is the Fermi level of the metal, and we use the boundary condition $\phi(\mathbf{r}, z = 0) = 0$ In both cases we shall work at a given $\mu_2 = \text{const.}$ with being its conjugated variable. The difference between the two types of junctions will reveal itself later when inhomogeneous fluctuations are treated.

It is anticipated that the system will have two types of states corresponding to the Bose-liquid homogeneous in the plane and to the plane Wigner crystal. Since both of these states break the continuous symmetry group and the system is effectively two-dimensional then there is no long-range ordering with $\langle \psi(\mathbf{R}) \rangle \neq 0$. Still in the both cases we can consider the states with an infinite correlation length of the BKT type. As for superfluidity these states are almost indistinguishable from three-dimensional systems with a long-range order.

We shall be able to consider part of the phase diagram where fluctuations of $\psi(\mathbf{R})$ and $\phi(\mathbf{R})$ are small

$$2e\phi(\mathbf{R}) - \mu = V(z) + v(\mathbf{R}, \tau); \qquad v \ll V$$
$$\psi(\mathbf{R}, \tau) = \rho^{1/2} \exp\{i\theta(\mathbf{R}, \tau)\}$$

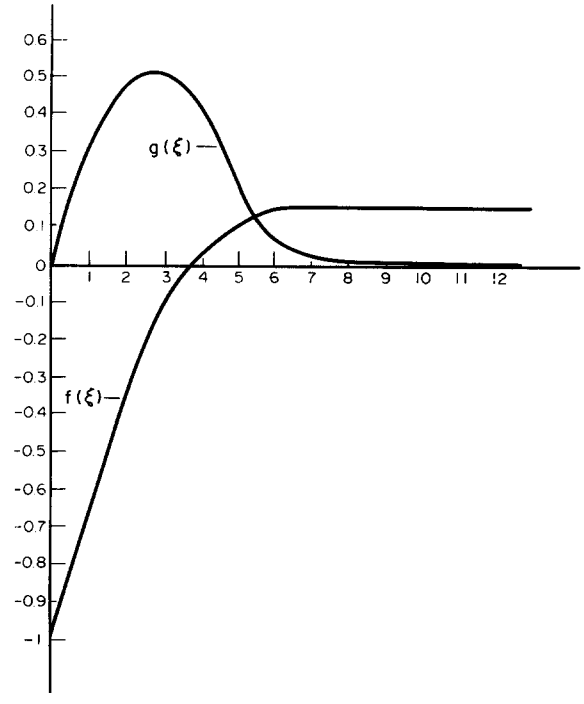


Fig. 1.

$$\rho = \rho_0(z) + \pi(\mathbf{R}, \tau) \qquad \pi \ll \rho_0$$
$$\rho_0 = \psi_0^2(z). \qquad (2)$$

The potential $V(z)$ and the wave function $\psi_0(z)$ must give the extremum of equation (1). So we find the selfconsistent system of Poisson and Schrödinger

$$\frac{1}{2M_\perp} \frac{d^2\psi_0(z)}{dz^2} + V(z)\psi_0(z) = 0$$
$$\frac{\epsilon_\perp}{4\pi} \frac{d^2V(z)}{dz^2} + 4e^2\psi_0^2(z) = 0. \qquad (3)$$

The solution of equation (3) is

$$V(z) = \mu f(\xi); \qquad \psi_0(z) = \sigma^{1/2} g(\xi) \int_0^\infty g^2(\xi) d\xi = 1$$
$$f(0) = -1; \qquad f(+\infty) = 0.16; \qquad (4)$$
$$\xi = z/d$$
$$d = (2M_\perp \mu)^{-1/2}; \qquad \sigma = \alpha \frac{r_\perp}{d^3}, \qquad r_\perp^2 = \hbar^2 \epsilon_1 / M_\perp \epsilon$$
$$\alpha = \frac{1'}{2^{7/2}\pi} \left. \frac{df}{d\xi} \right|_{\xi = 0} = 0.6. \qquad (5)$$

The functions $f(\xi)$ and $g(\xi)$ are plotted in Figure 1. They show a formation of the selfconsistent potential wall with the single localised eigenstate at the given eigenvalue μ. All bosons are condensed at this state. The relation

between σ and μ follows from equation (5). The asymptotical value of the BPs' electrochemical potential $\tilde{V}_\infty = V(\infty) = 0.16\mu > 0$ is higher than the edges for both the EC and IC spectra so that the BP volume concentration is zero. If the finite concentration of IC in the volume is taken into account a usual screening length r_D appears in the theory, then at $z \sim r_D$ the crossover takes to $V_\infty > \tilde{V}_\infty$ place. In the absence of volume doping we have evidently $V_\infty = 2e\phi(\infty) = E_2$. We assume here that

$$r_D \gg d \gg \xi_0, \tag{6}$$

where ξ_0 is the BPs' size. The crossover from $\tilde{V}_\infty$ to V_∞ at $z \sim r_D$ is described by the classical Debye screening whereas the crossover from μ to $\tilde{V}_\infty$ in the layer $z \sim d$ is a purely quantum effect.

3. FLUCTUATIONS AND ELECTROMAGNETIC PROPERTIES

In this section we shall consider small deviations of ψ and ϕ from the extremal solution (4, 5). Unlike the extremal solution the character of perturbations depends essentially on the boundary condition for ϕ, i.e. on the type of the junction. We shall consider the two limiting cases when the dielectric constant ϵ^- for a substance at $z < 0$ differs strongly from $\epsilon_\parallel$, $\epsilon_\perp$. Then it can be taken into account by means of the image charges $2e'$ which have been already included in (2).

1. Contact with a "metal": $\epsilon^- \gg \epsilon_\parallel, \epsilon_\perp, e' = -e$.
2. Contact with a "vacuum": $\epsilon^- \ll \epsilon_\parallel, \epsilon_\perp, e' = e$.

According to equations (1) and (2) we describe the small perturbations by the Lagrangian

$$\begin{aligned}\mathcal{L} = \sum_{k_\parallel,\omega}\Bigg[\int_0^\infty dz \Bigg\{ & i\omega\pi\theta - \frac{1}{2M_\perp}\rho_0(z)\left|\frac{\partial\theta}{\partial z}\right|^2 \\ & - \frac{1}{2M_\perp}\rho_0 k_\parallel^2|\theta|^2 - \frac{1}{8M_\parallel}\frac{k_\parallel^2|\pi|^2}{\rho_0(z)} - \frac{1}{8M_\perp}\left[\left|\frac{d\pi}{dz}\right|^2 \cdot \frac{1}{\rho_0} \right. \\ & \left. + \frac{d^2\rho_0}{dz^2}\cdot\frac{|\pi|^2}{\rho_0}\right]\Bigg\} \\ & - \frac{2e^2}{(\epsilon_\parallel\epsilon_\perp)^{1/2}}\int dz_1\, dz_2\, \pi(kz_1)\pi(-k, z_2)\cdot \\ & \frac{\exp\{-\nu|k_\parallel||z_1 - z_2|\} \mp \exp\{-\nu|k_\parallel||z_1 + z_2|\}}{|k_\parallel|}\end{aligned} \tag{7}$$

$$\nu = (\epsilon_\parallel/\epsilon_\perp)^{1/2}; \qquad k = (k_\parallel, \omega);$$

$$-k = (-k_\parallel, \omega); \qquad \pi = \pi(k, z); \qquad \theta = \theta(k, z),$$

where $\theta(k, z)$, $\pi(k, z)$ are the Fourier representations for $\theta(\tau, \mathbf{r}, z)$, $\pi(\tau, \mathbf{r}, z)$ in equation (2). The signs $\mp$ in the last term of equation (7) correspond to the cases 1 and 2.

Let us consider the spectrum of small oscillations. The inhomogeneity along z-direction does not allow to the exact diagonalization of equation (7). Nevertheless the long wavelength limit can be obtained with the help of themodynamical relations. We shall assume that all variables change slowly at the scales d, l so that they are quasistationary.

Let us introduce:

$$\begin{aligned}\delta(\mathbf{r}) &= \int_0^\infty dz\pi(\mathbf{r}, z)dz \\ \theta(\mathbf{r}) &= \theta(\mathbf{r}, z)|_{z=0}.\end{aligned} \tag{8}$$

In term of this variables we can write down the effective Lagrangian:

$$\mathcal{L} = \int d^2\mathbf{r}\left[\delta\frac{\partial\theta}{\partial\tau} + C\frac{\partial}{2M_\parallel}\frac{\partial\theta}{\partial\mathbf{r}}^2 + \frac{1}{2}\frac{\partial\mu}{\partial\sigma}\delta^2 + \mathcal{H}_c\right]$$

$$C \sim 1; \qquad \frac{d\mu}{d\sigma} = \frac{2}{3}\frac{\mu}{\sigma}. \tag{9}$$

The Coulomb energy $\mathcal{H}_c$ in equation (9) for the cases 1 and 2. becomes

$$\mathcal{H}_c = 2\frac{e^2}{\epsilon_\perp}\sum_{k_\parallel}|\delta(k_\parallel)|^2\begin{cases} c_1 d; \quad c_1 \sim 1 & \text{case 1.} \\ \left(\dfrac{\epsilon_\perp}{\epsilon_\parallel}\right)^{1/2}\dfrac{2}{k_\parallel} & \text{case 2.}\end{cases} \tag{10}$$

For the case 1. $\mathcal{H}_c$ reduces to the correction for $d\mu/d\sigma$ in equation (9) which happens to be of the same order and sign.

From equations (9) and (10) we find matrix of correlation function:

$$\begin{aligned}\hat{F}_{1,2}(\omega, k) &= \begin{pmatrix}\langle\theta(k)\theta(-k)\rangle & \langle\theta(k)\delta(-k)\rangle \\ \langle\delta(k)\theta(-k)\rangle & \langle\delta(k)\delta(-k)\rangle\end{pmatrix} \\ &= [\omega_n^2 + \omega_{1,2}^2(k)]^{-1}\begin{pmatrix}A_{1,2} & -\omega_n \\ \omega_n & B_{1,2}\end{pmatrix},\end{aligned} \tag{11}$$

where

$$\begin{aligned}&\omega_1(k) = uk; \qquad u^2 \sim \mu/M_\parallel; \qquad A_1 = \sigma\frac{\partial\mu}{\partial\sigma}; \\ &B_1 = \frac{\sigma k^2}{2M_\parallel} \\ &\omega_2(k) = \frac{\mu}{M_\parallel}k^2 + \frac{4\pi e^2\sigma}{M_\parallel(\epsilon_\parallel\epsilon_\perp)^{1/2}}k;\end{aligned} \tag{12}$$

$$A_2 = \frac{2\pi e^2}{\epsilon_\perp k}, \qquad B_2 = \frac{\sigma k^2}{2M_\parallel}, \tag{13}$$

$\omega_2(k)$ is a typical spectrum of the surface plasmons.

Formulas (13)–(15) are applicable at $|k| \ll k_m$ where the upper boundary k_m is determined by the crossover $uk_m \sim k_m^2/M_\parallel$, i.e. $k_m \sim d^{-1}(M_\parallel/M_\perp)^{1/2}$.

We can also estimate the fluctuations near the solutions (4, 5). With the help of equations (13)–(15) for the both cases 1 and 2. They are small provided that $l \ll r_\perp$, $T \ll \epsilon_\perp = \hbar^2/M_\perp r_\perp^2$. We can also use equations (13)–(15) to find the BPs' correlation function at large distances.

$$G(\mathbf{R}_1, \mathbf{R}_2) = \langle \psi(\mathbf{R}_1)\psi^+(\mathbf{R}_2)\rangle$$

$$\sim \psi_0(z_1)\psi_0^*(z_2)\left(\frac{d}{r}\right)^\gamma; \qquad \gamma = \frac{M_\parallel}{\rho}T, \tag{14}$$

as it should be for the two-dimensional Bose-liquid.

Let us consider the junction in the presence of the magnetic field described by the vector potential $\mathbf{A}$. We can neglect variation of $\mathbf{A}$ over the scale $z \sim d$. Then we use the free field Lagrangian for $\mathbf{A}$ in the whole volume together with the Lagrangian (9) taken after the gauge transformation

$$\nabla\theta \to \nabla\theta - 2\frac{e}{c}\mathbf{A}_0$$

$$A_0 = A_0(\mathbf{r}, \tau) = A(\mathbf{r}, z, \tau)|_{z=0}.$$

Similarly to [13] we come to the equation

$$\Delta\mathbf{A} = 8\pi\frac{e\sigma}{cM_\parallel}\left(\nabla\theta - \frac{2e}{c}\mathbf{A}\right)\delta(z). \tag{15}$$

For the Fourier transform of $A_0(\mathbf{r}, \tau)$ we find

$$A_0(\mathbf{k}, \omega) = 2\pi\frac{e\sigma}{mc}\frac{\mathbf{k}\theta\mathbf{k}}{|k| + \Lambda^{-1}}, \qquad \Lambda = \frac{mc^2}{2\pi\rho e^2}. \tag{16}$$

Λ in (16) is the London penetration length. Following the theory of superconductivity we also introduce the parameter: $\kappa = \Lambda/d \sim (\hbar c/e^2)^2 d^2/r_\perp r_\parallel$. Bearing in mind that $d^2/r_\perp r_\parallel$ cannot be too small ($d \gg \xi_0$) we conclude that $\kappa \gg 1$ so that the system would belong to the type II superconductor. The vortices in this system have been studied in [13].

4. THE PHASE DIAGRAM

First we consider qualitively finite temperatures effects. At low temperatures the system has the BKT type state characterised by the power law (14) and by the superfluid density $\rho_s = \rho_s(T)$. We see from equation (9) that at $T \to 0$ $\rho_s = \hbar^2\sigma/M_\parallel$. There is a critical temperature T_c such that $\rho_s(T_c) = 0$. Since we do not take into account the vertices we can find only a non-renormalised value $\rho_s^*(T)$ such that $\rho_s^*(T_c) \sim T_c$. There are two reasons for the decrease of ρ_s^* and increasing T.

i. Taking into account the creation of vortex pairs we obtain the usual estimate $T_c \leqslant \rho_s(0)$.

ii. The peculiar effect of our model are: the decrease of the Debye length, the population of higher localised states, etc. At $T \sim E_\perp$ the extremal solution (4) will be evidently smeared out, the Debye screening will dominate and $\rho_s^* < T$, $\rho_s = 0$ so that $T_c < E_\perp \sim \hbar^2/M_\perp r_\perp^2$.

Now we consider the competition of other phases at $T = 0$. At low σ near the injection threshold for BP the state of BP will depend on the type of the junction. For the contact with a "vacuum" we can anticipate formation of the Wigner crystal of the BPs which will dominate over their Bose-condensation at $T = 0$. Note that according to various calculations (see [14]) the region for the $2d$ – Wigner crystal existence is numerically narrow – $l/r > 100$. Then in the region $l \gtrsim r_\parallel, r_\perp$ there is still the homogeneous Bose-liquid and the superconductivity is possible like it was in the dense region $l < d(r_\parallel/r_\perp)^{1/2}$ which we have studied here.

Note also that according to [15] the Bose-condensation and the crystallisation generally do not exclude each other.

For the contact with a metals at low σ every BP will be bound to its image charge with the binding energy $\epsilon_{\text{dip}} \sim E_\perp$. The system may be considered as a dilute gas of neutral dipoles. Their concentration σ_d is finite above the threshold: $2\mu_0 > E_2 - \epsilon_{\text{dip}}$, i.e. $\mu > -\epsilon_{\text{dip}}$. The dipole–dipole interaction in the $2d$-system is start-ranged unlike the $3d$ case. Then we can describe it by the standard Lagrangian for the weakly nonlinear Bose-field $\varphi(\mathbf{r})$ of dipoles:

$$\mathcal{L} = i\varphi^*\frac{\partial\varphi}{\partial t} + \frac{1}{2M_\parallel}\left(\frac{\partial\varphi}{\partial\mathbf{r}}\right)^2 + (\mu - \epsilon_{\text{dip}})|\varphi|^2 + \lambda|\varphi|^4$$

$$\lambda = \int d^2\mathbf{r} v_{\text{dip}}(\mathbf{r}) \sim \int \frac{d^2\mathbf{r} e^2 \mathbf{r}_\perp^2}{\mathbf{r}^3} \sim e^2\mathbf{r}_\perp.$$

The Bose-condensation in such a $2d$ system has been as considered in [16]; At $T < \rho_s$, the superfluid density $\rho_s \sim \lambda/(\mu - \epsilon_{\text{dip}})$ emerges at $T < \rho_s$. For the case of the direct contact with a metal the dipole superfluidity will not manifest itself as superconductivity. Nevertheless our consideration is applicable also for the contact with a substance with the high dielectric constant: $\epsilon^- \gg \epsilon_\parallel, \epsilon_\perp$. Then $e' \neq -e$ and the dipole will have a small but finite charge $g = 2(e - e') \sim e^{\epsilon_\parallel \cdot \epsilon_\perp}/\epsilon$. Now the region for the W.C. $l \gg r_\parallel e^2/g^2$ will appear again. In any case at higher μ and σ the quantum screening will appear in the form as it has been discussed in Section 2.

5. CONCLUSIONS

Now we shall discuss some problems which can arise at an attempt to realise experimentally these phenomena. Evident restrictions for the electric fields at the contact with a dielectric $E < E_{max}$ require a rather low density of BP: $\sigma < E_{max}/4\pi$. To weaken the interaction with imperfections of the surface the BPs' was function should be maximally extended over z: $d \gg a$. According to (2) this condition also restricts σ from above. At the same time at low σ we meet with either the Wigner crystallization of the formation of dipoles (Section 4). The region for superconductivity will be larger for materials with high $\epsilon_{\|}$, $\epsilon_{\perp}$ and relatively low $M_{\|}, M_{\perp}$. The narrow gap dielectrics with BPs of large or moderate sizes when $M_{\|}$, $M_{\perp} \sim m^* \epsilon_2^2/\omega_{ph}^2$, where m^* is the effective mass and ω_{ph} is the phonon frequency, are interesting in this respect. These constraints except the one for the BPs. mass M_t for their motion between the chains are fulfilled in polymers, where generally

$$M_t = A(a_\perp^2 t_\perp)^{-1}; \qquad A = \exp(c\epsilon_2/\omega_{ph}^{-1})$$

$t_\perp$ is an electron interchain transfer integral, $a_\perp$ is an interchain spacing, $c \sim 1$. For the polaron in polyacetylene the Frank–Condon factor has been found to be $A = 10$ [17].

For this reason we would like to draw the attention to materials like polypyrrole and polythiophene which are exceptional due to their strong interchain overlap ($a_\perp = 3.6$ A, $t_\perp \sim 1$ eV). The situation is improved if the BP is near its stability threshold ($\epsilon_2 \lesssim 2\epsilon_1$) when $\Delta \gg \epsilon_2 \sim \omega_{ph}$ so that $M_{\|,\perp} \sim m^*$. This probably happens in Ti_4O_7 where according to the phase transition with the dissociation of BP at $T = 140$ K we have $\epsilon_2 \sim 10^2$ K $\sim \omega_{ph}$.

REFERENCES

1. C. Schlenker, S. Ahmed, R. Buder & M. Gourmala, *J. Phys. C*, **12**, 3503 (1979).
2. O.F. Schirmer & E. Salje, *J. Phys. C*, **13**, L1667 (1982).
3. L.W. Schacklette, R.R. Chance, D.M. Ivory, G.G. Miller & R.H. Bangman, *Synth. Met.* **1**, 307 (1979).
4. G.B. Street, T.C. Clarke, R.H. Geiss, V.Y. Lee, A. Nazzal, P. Pflugger & I.C. Scott, *J. Phys. Colloque* C3, **44**, 599 (1983).
5. T-C. Chung, J.H. Kaufman, A.J. Heeger & F. Wudl, *Phys. Rev.* **B30**, 702 (1984).
6. S. Brazovskii, **78**, 677 (1980); T. Bohr & S. Brazovskii. *J. Phys. C*, **16**, 1189 (1983).
7. S. Brazovskii, N. Kirova & V. Yakovenko, *J. Phys. Colloque*, C3, **44**, 1525 (1983).
8. Brazovskii & N. Kirova, *Pis'ma ZhE hsp Teor. Fiz.* **31**, 399 (1980); S. Brazovskii, N. Kirova & S. Matveenko, *Zh. Eksp. Teor. Fiz.* **86**, 743 (1984).
9. J.L. Brédas, B. Thémans, J.M. André, R.R. Chance, D.S. Boudreaux & R. Silbey. *J. Phys. Colloque*, C3, **44**, 373 (1983).
10. S. Matveenko, *Zh. Eksp. Teor. Fiz.* **86**, 1742, 1984.
11. N.H. March, W.H. Yang & S. Sampanhar, *The Manybody Problem in Quantum Mechanics.* Kembridge University Press (1967).
12. A. Alexandrov & J. Runninger, *Phys. Rev. B*, **24**, 1164 (1981).
13. J. Perl, *Appl. Phys. Lett.*, **5**, 65 (1964).
14. R.W. Hockey & T.R. Brown, *J. Phys. C*, **8**, 1813 (1975), C.C. Grimes & G. Adams, *Phys. Rev. Lett.* **42**, 795 (1979).
15. I. Dzjaloshinskii & P. Kondratenko, *Zh. Eksp. Teor. Phys.* **77**, 241 (1975).
16. V. Popov, *Teor. Mat. Fiz.* **27**, 745 (1979).
17. S. Jeyadev & J.R. Schrieffer, *Phys. Rev. B*, **30**, 3620 (1984).