Review Article

Competition of Superconductivity and Charge Density Waves in Cuprates: Recent Evidence and Interpretation

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Explicit and implicit experimental evidence for charge density wave (CDW) presence in high- T_c superconducting oxides is analyzed. The theory of CDW superconductors is presented. It is shown that the observed pseudogaps and dip-hump structures in tunnel and photoemission spectra are manifestations of the same CDW gapping of the quasiparticle density of states. Huge pseudogaps are transformed into modest dip-hump structures at low temperatures, T , when the electron spectrum superconducting gapping dominates. Heat capacity jumps at the superconducting critical temperature and the paramagnetic limit are calculated for CDW superconductors. For a certain range of parameters, the CDW state in a d -wave superconductor becomes reentrant with T , the main control quantity being a portion of dielectrically gapped Fermi surface. It is shown that in the weak-coupling approximation, the ratio between the superconducting gap at zero temperature $\Delta(T = 0)$ and T_c has the Bardeen-Cooper-Schrieffer value for s -wave Cooper pairing and exceeds the corresponding value for d -wave pairing of CDW superconductors. Thus, large experimentally found values $2\Delta(T = 0)/T_c \approx 5 \div 8$ are easily reproduced with reasonable input parameter values of the model. The conclusion is made that CDWs play a significant role in cuprate superconductivity.

1. Introduction

Ever since the earliest manifestations of high- T_c superconductivity were found in 1986 [1], the whole theoretical power [2–22] has been applied to explain and describe various normal and superconducting properties of various oxide families with critical temperatures, T_c , ranging up to 138 K to date [23–27]. Unfortunately, even conceptual understanding of the mechanisms and character of superconductivity in cuprates is still lacking. Strictly speaking, there is a number of competing paradigms, every of them pretending to be “the theory of superconductivity” (see, e.g., [2]) but not recognized as such by other respected experts in the field.

After the discovery of high- T_c oxides, experimentalists found several other superconducting families with T_c higher than 23.2 K reached by the precuprate record-holder, Nb_3Ge [28, 29]. For instance, one may refer to fullerenes [30, 31],

doped bismuthates [32–34], hafnium nitrides [35, 36], and magnesium diborides [37–40]. One should also mention more controversial cases of superconducting oxides H_xWO_3 with $T_c \approx 120$ K [41] and $\text{Sr}_{0.9}\text{La}_{0.1}\text{PbO}_{3-\delta}$ with $T_c \approx 65$ K [42]. Finally, an unexpected and counter-intuitive discovery of the iron-based oxypnictide [43, 44] or oxygen-free pnictide [45] layered superconductors with T_c over 50 K has been made recently (see also reviews [46–49]).

Presumably, the latter materials with FeAs layers have been overlooked as possible candidates for high- T_c superconductors, since Fe ions in solids usually possess magnetic moments, which promote magnetic ordering, the latter being detrimental to superconductivity, especially the spin-singlet one [50–55]. Strictly speaking, such an omission is of no surprise because superconductivity in oxides is rather gentle, sensible to impurities, including the excess or deficiency of oxygen [56] in these nonstoichiometric [57, 58]

compounds. Recent discovery [59] of previously unnoticed high- T_c superconductivity in parent compounds $T'-R_2\text{CuO}_4$ ($R = \text{Pr, Nd, Sm, Eu, Gd}$) is very symptomatic in this regard, since an accurate removal of apical oxygen from thin films raised T_c from exact zero (those compositions were earlier considered by theoreticians as typical correlated Mott-Hubbard insulators) to 32.5 K for Nd_2CuO_4 . As for the ferroarsenide family, one of its members, $\text{EuFe}_2(\text{As}_{0.7}\text{P}_{0.3})_2$, reveals a true superconducting transition at 26 K, followed by the ferromagnetic ordering of Eu^{2+} magnetic moments below 20 K, coexisting with superconductivity [60], which is quite unusual in view of the antagonism indicated above between two kinds of cooperative phenomena.

What is more, none of the mentioned superconductors except $\text{Ba}_{1-x}\text{K}_x\text{BiO}_3$ [33] were discovered due to theoretical predictions. Hence, one may consider the theoretical discovery of $\text{Ba}_{1-x}\text{K}_x\text{BiO}_3$ as an accidental case, since, according to the well-known chemist Cava: “one of the joys of solid state chemistry is its unpredictability” [61]. The same opinion was expressed by the other successive chemist Hosono: “understanding the mechanism with respect to predicting the critical temperature of a material is far from complete at the present stage even for brilliant physicists. Such a situation provides a large opportunity including a good luck for material scientists who continue the exploration for a new material, not limited to superconductors, and a new functionality based on their own view points” [48]. That is why Pickett recently made a sad remark that “the next breakthrough in superconductivity will not be the result of surveying the history of past breakthroughs” [62]. It means that *microscopic* theories of superconductivity are incapable of describing specific materials precisely, although *together* they give an adequate overall picture. In this connection, the failure of the most sophisticated approaches to make *any* prediction of true or, at least “bare” T_c , (provided that the corresponding T_c -value is not known *a priori*) despite hundreds of existing superconductors with varying fascinating properties, forced Phillips [63] to reject all apparently *first-principle* continuum theories in favor of his own percolative filamentary theory of superconductivity [64–67] (see also the random attractive Hubbard model studies of superconductivity [68, 69] and the analysis of competition between superconductivity and charge density waves studied in the framework of similar scenarios [70–72]). We totally agree with such considerations in the sense of the important role of disorder in superconductors with high T_c on the verge of crystal lattice instability [73–83]. Nevertheless, it is questionable whether a simple one-parameter “master function” of [63, 67] would be able to make *quantitative* and practically precise predictions of T_c . As for the qualitative correctness of the dependence T_c versus weighted number $\langle R \rangle$ of Pauling resonating valence bonds [63, 67, 84], it can be considered at least as a useful guideline in the superconductivity ocean. The phenomenological character of the master function (chemical trend diagram) $T_c(\langle R \rangle)$ is an advantage rather than a shortcoming of this approach, as often happens in the physics of superconductors (see, e.g., more or less successful criteria of superconductivity with different extent of phenomenology [85–98]).

On the other hand, attempts to build sophisticated microscopic theories of the boson-mediated Bardeen-Cooper-Schrieffer (BCS) attraction, treating the Coulomb repulsion as a single Coulomb pseudopotential constant μ^* , are incapable of predicting actual critical superconducting properties [63, 91, 99–101]. The same can be said [67] about Hubbard-Hamiltonian models with extremely strong repulsive Coulomb energy parameter U , which is formally based on the opposite ideology (see, e.g., [102, 103]). As an example of the theories described above, one can indicate work [104], where the strong-coupling Eliashberg equations for the electron-phonon mechanism of superconductivity [105, 106] were solved numerically taking into account even vertex corrections and treating the dispersive Coulomb interaction not on equal footing, but as a simple constant μ^* . In this connection, it seems that the prediction of [104] that the maximal T_c for new iron-based superconductors is close to 90 K is unjustified. Of course, the same is true for other studies of such a kind.

It is remarkable that, for hole- and electron-doped cuprates, there is still no clarity concerning the specific mechanisms of superconductivity [17, 107–115] and the order parameter symmetry [109, 116–130], contrary to the “official” viewpoint [131–133]) and even the very character of the phenomenon (in particular, there have been furious debates concerning the Cooper pairing versus boson condensation dilemma in cuprates [8, 134, 135]). The same seems to be true for other old and new “exotic” superconductors [46, 107, 108, 136–153], their exoticism being in essence a degree of our ignorance.

It would be of benefit to consider all indicated problems in detail for all classes of superconductors and show possible solutions. Unfortunately, it cannot be done even in the scope of huge treatises (see, e.g., [154–157]). The objective of this review is much more modest. Specifically, it deals mostly with high- T_c cuprate materials, other superconductors being mentioned only for comparison. Moreover, in the present state of affairs, it would be too presumptuous to pretend to cover all aspects of the oxide superconductivity. Hence, we will restrict ourselves to the analysis of lattice instabilities and concomitant charge density waves (CDWs) in high- T_c oxides. Their interplay with superconductivity is one of the fascinating and fundamentally important phenomena observed in cuprates and discussed by us earlier [158–160]. Nevertheless, in this rapidly developing branch of the solid state physics, many new theories and experimental data on various CDW superconductors appeared during last years. They are waiting for both unbiased and thorough analysis. This article discusses this new information, referring the reader to our previous reviews for more general and established issues, as well as some cumbersome technical details.

The outline of this review is as follows. In Section 2, for the sake of completeness, we briefly consider possible mechanisms of superconductivity in cuprates, the problem of the relationship between BCS pairing and Bose-Einstein condensation (BEC), and the multigapness of the superconducting order parameter. Section 3 is devoted to the experimental evidence for CDWs, the so-called

pseudogaps, dip-hump structures, and manifestations of intrinsic inhomogeneity in cuprate materials. The original theory of CDW superconductors and the interpretation of CDW-related phenomena in high- T_c oxides are presented in Section 4. At the end of Section 4, some recent data on coexistence between superconductivity and spin density waves (SDWs)—a close analogue of CDWs—are covered. This topic became hot once more after the discovery of ferropnictides [43–48, 161]. Short conclusions are made in Section 5.

2. Considerations on Peculiarities and Mechanisms of Superconductivity in Oxides

When $\text{BaPb}_{1-x}\text{Bi}_x\text{O}_{3-\delta}$ (BPB) was shown [162] to be a superconductor with a huge (at that time!) $T_c \approx 13$ K for $x \approx 0.25$, a rather low concomitant concentration of current carriers $n \approx 1.5 \div 4.5 \times 10^{21} \text{ cm}^{-3}$, and poor electric conductivity [56] (the phase diagram of BPB is extremely complex, with a number of partial metal-insulator structural transitions [56, 163–167]), it looked like an exception. Now, it is fully recognized that oxides with highest T_c are bad metals from the viewpoint of normal state conductivity [168]. In particular, the mean free path of current carriers is of the order of the crystal lattice constant, so that the Ioffe-Regel criterion of the metal-insulator transition [169] is violated. Moreover, there exists an oxide superconductor $\text{SrTiO}_{3-\delta}$ with a tiny maximal $T_c \approx 0.5$ K, attained by doping, but an extremely small $n < 10^{20} \text{ cm}^{-3}$ [170]. Note that the undoped semiconducting $\text{SrTiO}_{3-\delta}$ is so close to the metal-insulator border that it may be transformed into a metal by the electrostatic-field effect [171] (this technique has been successfully applied to other oxides [172]). Moreover, a two-dimensional metallic layer has been discovered [173] at the interface between two insulating oxides $\text{LaAlO}_{3-\delta}$ and $\text{SrTiO}_{3-\delta}$, which was later found to be superconducting with $T_c \approx 0.2$ K [174, 175]. The appearance of superconductivity at nonmetallic charge carrier densities in oxides of different classes comprises a hint that it is not wise to treat various oxide families separately (see, e.g., [176]), all of them having similar perovskite-like ion structures [23, 25, 26, 177–181] and similar normal and superconducting properties [27], whatever the values of their critical parameters are. As for the apparent dispersion of the latter among superconducting oxides, it mostly reflects their conventional exponential dependences on atomic and itinerant-electron characteristics [9, 10].

The junior member of the superconducting oxide family, $\text{SrTiO}_{3-\delta}$, demonstrates (although not in a spectacular manner) several important peculiarities, which are often considered as properties intrinsic primarily to high- T_c cuprates. Indeed, in addition to the low n , this polar, almost ferroelectric [182, 183], material was shown to reveal polaron conductivity [184] and is suspected to possess bipolaron superconductivity [185–187], first suggested by Vinetskii almost 50 years ago [188]. It means that $\text{SrTiO}_{3-\delta}$ might be not a Bardeen-Cooper-Schrieffer (BCS) superconductor [189] with a large coherence length $\xi_0 \gg a_0$, where a_0

is the crystal lattice constant, but most likely an example of a material with $\xi_0 \approx a_0$, so that a Bose-condensation of local electron pairs would occur at T_c , according to the Schafroth-Butler-Blatt scenario [190] or its later extensions [8, 134, 135, 191–199].

The concept of bipolarons (local charge carrier pairs) has been later applied to BPB [200–203], $\text{Ba}_{1-x}\text{K}_x\text{BiO}_{3-\delta}$ (BKB, $T_c \leq 30$ K [204, 205]) [203, 206, 207] and cuprates [195, 199, 208–211]. It was explicitly shown for BPB and BKB by X-ray absorption spectroscopy [203] that bipolaronic states and CDWs coexist and compete, which might lead, in particular, to the observed nonmonotonic dependence $T_c(x)$ [212]. At the same time, Hall measurements demonstrate that the more appropriate characteristics $T_c(n)$ is monotonic [56, 213, 214], so that the expected suppression of T_c at high n as a consequence of screening of the electron-phonon matrix elements [99, 215, 216] is not achieved here as opposed to the curve $T_c(n)$ [170] in reduced samples of $\text{SrTiO}_{3-\delta}$. As for cuprates, the bipolaron superconductivity mechanism, as well as any other BEC scheme, in its pure state would require an existence of the preformed electron (hole) pairs (bipolarons), which might be the case [177, 217], and a prior destruction of the Fermi surface (FS), the condition contradicting observations (see, e.g., [218]). Therefore, boson-fermion models for charge carriers in superconductors was introduced [134, 219–225] and, later on, severely criticized [226, 227]. In any case, the available objections concern the bipolaronic mechanism of superconductivity itself, the occurrence of polaronic effect in oxides with high dielectric permittivities raising no doubt [115, 177, 199, 228–232].

It is remarkable that the boson-fermion approach mentioned above is not a unique tool for describing superconductivity in complex systems. A necessary “degree of freedom” connected to another group of charge carriers has been introduced, for example, as the so-called $(-U)$ -centers [233–235], earlier suggested by Anderson [236] as a phenomenological reincarnation of bipolarons in amorphous materials [188]. Independently, narrow-band nondegenerate charge carriers submerged into the sea of itinerant electrons were proposed for cuprates as another, not fully hybridized kind of the “second heavy component” [237, 238]. For completeness, we should also mention a quite different model involving a second heavy charge carrier subsystem (d -electrons in transition metals [239] or heavy holes in degenerate semiconductors [240]), necessary to convert high-frequency Langmuir plasmons intrinsic to the itinerant electron component into the ion-acoustic collective excitation branch, in order that a high- T_c superconductivity would appear. Those hopes, however, lack support from any evidence in natural or artificial systems (see the analysis of plasmon mechanisms [206, 241–247], the optimism of some authors seems to us and others [248] a little bit exaggerated). As can be readily seen from the References given above, all nonconventional approaches, rejecting or generalizing the BCS scheme and going back to the explanations of a relatively weak superconductivity in degenerate semiconductors [138, 191, 215, 249–252], have been applied to every family of superconducting oxides, including cuprates.

Strontium titanate became a testing ground [253] of one further attractive idea (based on the same concept of several interacting charge carrier components) of two-gap or multigap superconductivity, with the interband interplay being crucial to the substantial increase of T_c and other critical parameters. The corresponding models came into being in connection with the transition s - d metals [254, 255]. They were subsequently applied to analyze superconductivity in multivalley semiconductors [256, 257], high- T_c oxides [231, 258–266], MgB_2 [40, 267–269], ZrB_{12} [270], V_3Si [271], $\text{Mg}_{10}\text{Ir}_{19}\text{B}_{16}$ [272], $\text{YNi}_2\text{B}_2\text{C}$ [273], NbSe_2 [274, 275], $\text{R}_2\text{Fe}_3\text{Si}_5$ ($R = \text{Lu, Sc}$) [276], $\text{Sc}_5\text{Ir}_4\text{Si}_{10}$ [277], $\text{Na}_{0.35}\text{CoO}_2 \cdot 1.3\text{H}_2\text{O}$ [278] as well as pnictides $\text{LaFeAsO}_{1-x}\text{F}_x$ [279], $\text{LaFeAsO}_{0.9}\text{F}_{0.1}$ [280], $\text{SmFeAsO}_{0.9}\text{F}_{0.1}$ [281], and $\text{Ba}_{0.55}\text{K}_{0.45}\text{Fe}_2\text{As}_2$ [282], $\text{Ba}_{1-x}\text{K}_x\text{Fe}_2\text{As}_2$ with $T_c \approx 32\text{ K}$ [283]. We did not explicitly include into the list such modifications of magnesium diboride as $\text{Mg}_{1-x}\text{Al}_x\text{B}_2$ or $\text{Mg}(\text{B}_{1-x}\text{C}_x)_2$, and so forth.

Since, instead of one, two or more well-separated superconducting energy gaps, a continuous, sometimes wide, gap distribution is often observed (see results for Nb_3Sn in [284] and MgB_2 in [285–289]), the original picture of the gap multiplicity in the momentum, $\mathbf{k}$, space loses its beauty, whereas the competing scenario [76, 290] of the spatial ($\mathbf{r}$ -space) extrinsic or intrinsic gap spread becomes more adequate and predictive [77–79]. For the case of cuprates, it has been recently shown experimentally that the spread is really spatial, but corresponds to the pseudogap (CDW gap) rather than its superconducting counterpart, the latter most probably being a single one [291] (see also the discussion in [83] and below).

In accordance with what was already mentioned, the application of very different, sometimes conflicting, models to oxide families, including cuprates, means an absence of a deep insight into the nature of their superconducting and normal state properties. We are not going to analyze here the successes and failures of the *microscopic* approaches to high- T_c superconductivity in detail; instead we want to emphasize that even the boson-mediators (we accept the applicability of the Cooper-pairing concept to oxides on the basis of crucial flux-quantization experiments [292, 293]) are not known for sure. Indeed, at the early stages of the high- T_c studies, magnons were considered as glue, coupling electrons or holes. The very temperature-composition (doping) phase diagrams supported this idea, since undoped and slightly doped oxides were found antiferromagnetic [26, 103, 294–304]. However, a plethora of theories suggesting virtual spin fluctuations as the origin of superconductivity in high- T_c oxides and leading to the $d_{x^2-y^2}$ symmetry of the superconducting order parameter have been developed [6, 11, 15, 103, 302, 305–310].

Fortunately for the scholars, it became clear that reality is richer for oxides than was expected, so that (i) the order parameter may include a substantial s -wave admixture [109, 116–129]; and (ii) phonons still exist in perovskite crystal lattices, inevitably affecting or, may be, even determining the pairing process [4, 10, 112, 115, 311], not to talk about polaron and bipolaron effects discussed above. It should be noted that there are reasonable scenarios of

d -wave order parameter symmetry in the framework of the electron-phonon interaction alone [208, 312–316] (a similar conclusion was made for the case of plasmon mechanism [317]).

At the same time, if one adopts a substantial (crucial?) role of spin-fluctuation mechanism in superconductivity, the ubiquitous phonons can (i) be neutral to the dominant d -wave pairing; (ii) act synergetically with spin fluctuations; (iii) or reduce T_c , as it would have been for switched-off phonons. The existing theories support all three variants, although some authors cautiously avoid any direct conclusions [103]. For instance, Kulić demonstrated the destructive interference between both mechanisms of superconductivity [4]. Phononic reduction of the magnetically induced T_c was claimed in [308, 318], whereas anisotropic phonons seem to enhance T_c , thus obtained [319]. Finally, according to [228, 320, 321], spins and phonons act constructively in cuprates. Once again, the microscopic approach was incapable of unambiguously predicting a result for the extremely complex system.

One should bear in mind that the problem is much wider than the interplay between spin excitations and phonons. Namely, it is more correct to consider the interplay between Coulomb inter-electron and electron-lattice interactions [232, 322]. Of course, the latter is also Coulombic in nature, phonons being simply an ion sound, that is, ion Langmuir plasma oscillations [323] screened in this case by degenerate light electrons [324] (thus, acoustic phonons constitute a similar phenomenon as the acoustic plasmons in the electron system [239, 240] with an accuracy to frequencies). One of the main difficulties is how to separate the metal constituents in order that some contributions would not be counted twice [322, 325–333]. Since it is possible to do rigorously only in primitive plasma-like models [91, 92, 99, 251, 322], the problem has not been solved. Therefore, empirical considerations remain the main source of future success for experimentalists, as it happened, for example, in the case of MgB_2 [37].

3. CDWs and CDW-Related Phenomena in Cuprates

The reasoning presented in Section 2 demonstrates that for the objects concerned, it is insufficient to rely only on microscopic theories, so that phenomenological approaches should deserve respect and attention. In actual truth, they might not be less helpful in understanding the normal and superconducting properties of cuprates, being generalizations of a great body of experimental evidence collected during last decades. In this section, we are going to show that two very important features are common to all high- T_c families. Specifically, these are the intrinsic inhomogeneity of nonstoichiometric superconducting ceramic and single crystalline samples [334–343] and the persistence of CDWs [340, 344, 345] and other phenomena, which we also consider as CDW manifestations (dip-hump structures, DHS [339, 346–348], and pseudogaps below and above T_c [349–358] in tunneling spectra and angle-resolved photoemission spectra, ARPES).

FIGURE 1: (Color online) Fermi surface nesting; and tight-binding-calculated Fermi surface (solid black curve) of optimally doped $\text{Bi}_2\text{Sr}_2\text{CuO}_{6+\delta}$ based on ARPES data [373]. The nesting wave vector (black arrow) in the antinodal flat band region has length $2\pi/6.2a_0$. Underdoped $\text{Bi}_2\text{Sr}_2\text{CuO}_{6+\delta}$ Fermi surfaces (shown schematically as red dashed lines) show a reduced volume and longer nesting wave vector, consistent with a CDW origin of the doping-dependent checkerboard pattern reported here (Taken from [344]).

CDWs were seen directly as periodic incommensurate structures in superconducting $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ (BSCCO) using various experimental methods [12, 334, 359–370]. Photoemission studies reveal the $4a_0 \times 4a_0$ charge-ordered “checkerboard” state in $\text{Ca}_{2-x}\text{Na}_x\text{CuO}_2\text{Cl}_2$ [371], and tunnel measurements visualized the same kind of ordering in BSCCO [370]. Scanning tunnel microscopy (STM) measurements found CDWs in $\text{Bi}_2\text{Sr}_{1.4}\text{La}_{0.6}\text{CuO}_{6+\delta}$ ($T_c^{\text{max}} \approx 29\text{ K}$) with an incommensurate period and CDW wave vectors $\mathbf{Q}$ depending on oxygen doping degree [340]. The same method revealed nondispersive (energy-independent) checkerboard CDWs in $\text{Bi}_{2-y}\text{Pb}_y\text{Sr}_{2-z}\text{La}_z\text{CuO}_{6+x}$ ($T_c \approx 35\text{ K}$ for the optimally doped composition) [344]. In this case, $\mathbf{Q}$ substantially depends on doping, rising from $\pi a_0^{-1}/6.2$ in an optimally doped sample to $\pi a_0^{-1}/4.5$ for an underdoped sample with $T_c \approx 25\text{ K}$. It is easily explained by the authors taking into account the shrinkage of the hole FS with decreasing hole number, so that the vector $\mathbf{Q}$ that links the flat nested FS sections grows, whereas the CDW period decreases (see Figure 1). One should note that, in the presence of impurities (e.g., an inevitably non-homogeneous distribution of oxygen atoms), the attribution of the observed charge order (if any) to unidirectional or checkerboard type might be ambiguous [372].

A similar coexistence of CDWs and superconductivity was observed in a good many different kinds of materials with a reduced dimensionality of their electron system, so that the corresponding FS includes nested (congruent) sections [158–160]. For completeness, we will add some new cases discovered after our previous reviews were published. First of all, the analogy between CDWs in cuprates and layered dichalcogenides was proved by ARPES [352, 374–376]. It should be noted that CDW competition with

superconductivity in cuprates was supposed as early as in 1987 on the basis of heat capacity and optical studies [377], whereas the similarity between high- T_c oxides and dichalcogenides was first noticed by Klemm [378, 379]. Additionally, a new dichalcogenide system Cu_xTiSe_2 was found with coexisting superconductivity and CDWs (at $0.04 < x < 0.06$) [380, 381]. The coexistence between two phenomena was observed in the organic material $\alpha\text{-(BEDT-TTF)}_2\text{KHg(SCN)}_4$, but superconductivity was attributed to boundaries between CDW domains, where the CDW order parameter is suppressed [382]. High-pressure studies of another organic conductor $(\text{Per})_2[\text{Au(mnt)}_2]$ revealed an appearance of superconductivity after the CDW suppression [383]. Still, it remained unclear, whether some remnants of CDWs survived in the superconducting region of the phase diagram. Application of high pressure also suppressed CDWs in the compound TbTe_3 at about $P = 2.3\text{ GPa}$, inducing superconductivity with $T_c \approx 1.2\text{ K}$, enhanced to 4 K at $P = 12.4\text{ GPa}$ [384], the behavior demonstrating the competition of Cooper and electron-hole pairings for the FS [385, 386]. The same experiments in this quasi-two-dimensional material revealed two kinds of CDW anomalies merging at $P = 2.3\text{ GPa}$, as well as antiferromagnetism, which makes this object especially promising. Finally, CDWs were found in another superconducting oxide $\text{Na}_{0.3}\text{CoO}_2 \cdot 1.3\text{ H}_2\text{O}$ by specific heat investigations [387–389], showing two-energy-gap superconductivity for as-prepared samples and non-superconducting CDW dielectrized state after ageing of the order of days. The sample ageing is a situation widely met for superconductors [390, 391], whereas the dielectrization of as-synthesized superconducting ceramic samples accompanied by a transformation of bulk superconductivity into a percolating one with the CDW background was observed for BPB long ago [56, 392]. Nevertheless, such a scenario was not proved directly at that time, while the bulk heat capacity peak in $\text{Na}_{0.3}\text{CoO}_2 \cdot 1.3\text{ H}_2\text{O}$ [387–389] unequivocally shows the emergence of CDWs instead of superconductivity.

We emphasize that CDWs compete with superconductivity, whenever they meet on the same FS. This is the experimental fact, which agrees qualitatively with a number of theories [385, 386, 393–397].

Returning to cuprates, we want to emphasize that the existence of pseudogaps above and below T_c is one of their most important features. Pseudogap manifestations are diverse, but their common origin consists in the (actually, observed) depletion of the electron densities of states (DOS). It is natural that tunnel and ARPES experiments, which are very sensitive to DOS variations, made the largest contribution to the cuprate pseudogap data base (see references in our works [81–83, 158–160]). Recent results show that the concept of two gaps (the superconducting gap and the pseudogap, the latter considered here as a CDW gap) [82, 352, 353, 357, 377, 398–404] begins to dominate in the literature over the one-gap concept [211, 355, 405–416], according to which the pseudogap phenomenon is most frequently treated as a precursor of superconductivity (for instance, a gas of bipolarons that Bose-condenses below T_c [413] or a d -wave superconducting-like state without a long-range phase rigidity [416]). The main arguments,

which make the one-gap viewpoint less probable, is the coexistence of both gaps below T_c [349, 417], their different position in the momentum space of the two-dimensional Brillouin zone [351, 353, 356, 418, 419], and their different behaviors in the external magnetic fields $\mathbf{H}$ [420], for various dopings [417], and under the effects of disordering [419].

Nevertheless, some puzzles still remain unresolved in the pseudogap physics. For instance, Kordyuk et al. [352] found that the pseudogap in $\text{Bi(Pb)}_2\text{Sr}_2\text{Ca(Tb)Cu}_2\text{O}_{8+\delta}$ revealed by ARPES is nonmonotonic in T . Such a behavior, as they indicated, might be related to the existence of commensurate and incommensurate CDW gaps, in a close analogy with the case of dichalcogenides [421]. Another photoemission study of $\text{La}_{1.875}\text{Ba}_{0.125}\text{CuO}_4$ has shown [354] that there seems to be two different pseudogaps: a d -wave-like pseudogap—a precursor to superconductivity—near the node of the truly superconducting gap and a pseudogap in the antinodal momentum region—it became more or less familiar to the community during last years [350, 351, 353, 356, 403, 418, 419] and is identified by us as the CDW gap.

Despite existing ambiguities, the most probable scenario of the competition between CDW gaps (pseudogaps) and superconducting gaps in high- T_c oxides, in particular, in BSCCO, includes the former emerging at antinodal (nested) sections of the FS and the latter dominating over the nodal sections (see Figure 2, reproduced from [403], where BSCCO was investigated, and results for $(\text{Bi,Pb})_2(\text{Sr,L a})_2\text{CuO}_{6+\delta}$ presented in [356]). Since CDW gaps are much larger than their superconducting counterparts, the simultaneous existence of the superconducting gaps in the antinodal region might be overlooked in the experiments. This picture means that the theoretical model of the partial dielectric gapping (of CDW origin or caused by a related phenomenon—spin density waves, SDWs) belonging to Bilbro and McMillan [385] (see also [56, 158–160, 386, 397, 422–428]) is adequate for cuprates. On the other hand, the coexistence of CDW and superconducting gaps, each of them spanning the whole FS [429–432], can happen only for extremely narrow parameter ranges [433]. Moreover, as is clearly seen from data presented in Figure 2 [403] and a lot of other measurements for different classes of superconductors, complete dielectric gapping has not been realized. The reason is obvious: nested FS sections cannot spread over the whole FS, since the actual crystal lattice is always three-dimensional and three-dimensionality effects lead to the inevitable FS warping detrimental to nesting conditions formulated below.

It is interesting that pseudogaps were also observed in oxypnictides $\text{LaFeAsO}_{1-x}\text{F}_x$ and $\text{LaFePO}_{1-x}\text{F}_x$ by ARPES [434] and $\text{SmFeAsO}_{0.8}\text{F}_{0.2}$ by femtosecond spectroscopy [435], where SDWs might play the same role as CDWs do in cuprates. At the same time, in iron arsenide $\text{Ba}_{1-x}\text{K}_x\text{Fe}_2\text{As}_2$, photoemission studies detected a peculiar electronic ordering with a $(\pi/a_0, \pi/a_0)$ wave vector [436], a true nature of which is still not known, but which might be related either to the magnetic reconstruction of the electron subsystem (SDWs) and/or to structural transitions (when CDWs accompanied by periodic crystal lattice distortions emerge in the itinerant electron liquid near the structural transition

FIGURE 2: (Color online) Schematic illustrations of the gap function evolution for three different doping levels of $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$. (a) Underdoped sample with $T_c = 75$ K. (b) Underdoped sample with $T_c = 92$ K. (c) Overdoped sample with $T_c = 86$ K. At 10 K above T_c there exists a gapless Fermi arc region near the node; a pseudogap has already fully developed near the antinodal region (red curves). With increasing doping, this gapless Fermi arc elongates (thick red curve on the Fermi surface), as the pseudogap effect weakens. At $T < T_c$ a d -wave like superconducting gap begins to open near the nodal region (green curves); however, the gap profile in the antinodal region deviates from the simple $d_{x^2-y^2}$ form. At a temperature well below T_c ($T \ll T_c$), the superconducting gap with the simple $d_{x^2-y^2}$ form eventually extends across entire Fermi surface (blue curves) in (b) and (c) but not in (a). (Taken from [403].)

temperature T_d [437, 438]). The interplay between structural and magnetic instabilities is important for pnictides [161],

since, for example, structural and SDW anomalies appear jointly at 140 K in BaFe_2As_2 [439]. It is not inconceivable that pnictides may be a playground for density waves as well as high- T_c oxides, with a rich variety of attendant manifestations.

The DHS is another visiting card of cuprates, being a peculiarity in tunnel and photoemission spectra at low $T \ll T_c$ and energies much higher than those of coherent superconducting peaks [81–83, 160, 339, 347, 348, 440, 441]. It is remarkable that in the S-I-N tunnel junctions, where S, I, and N stand for a high- T_c superconductor, an insulator, and a normal metal, respectively, a DHS might appear for either one bias voltage V polarity only [347] or both [442, 443], depending on the specific sample. In S-I-N junctions, current-voltage-characteristics (CVCs) with two symmetrically located DHSs (one per branch) are also observed, but with amplitudes that can differ drastically [442, 443]. In S-I-S symmetric junctions, DHS structures are observable (or not) in CVC branches of both polarities simultaneously [347], which seems quite natural. It is very important that although the CVC for every in the series of S-I-N junctions with BSCCO as superconducting electrodes was nonsymmetric, especially due to the presence of the DHS, the CVC obtained by averaging over an ensemble of such junctions turned out almost symmetric, or at least its nonsymmetry turned out much lower than the nonsymmetry of every CVC taken into consideration [443].

There is quite a number of interpretations concerning this phenomenon [347, 444–450]. We have discussed most of them in detail in our previous publications, whereas our theory and necessary reference to other models will be presented below.

STM mapping of high- T_c oxide samples revealed substantial inhomogeneities of energy gap spatial distribution [334, 336, 338, 339, 341–343, 363, 370, 441, 451–459]. The same conclusion was made from the interlayer tunneling spectroscopy [460, 461], more conventional S-I-N tunnel (point-contact) studies [440, 442], optical femtosecond relaxation spectroscopy [337], and inelastic neutron scattering measurements [335]. It is quite natural that some inhomogeneity should exist, since the oxygen content is always nonstoichiometric in those compounds [304]. Indeed, correlations were found between oxygen dopant atom positions and the nanoscale electronic disorder probed by STM [336]. The problem has been recently investigated theoretically making allowance for electrostatic modulations of various system parameters by impurity atoms [462].

Nevertheless, the gap distributions occurred to be anomalously large, with sometimes conspicuous two-peak structures in BSCCO [451, 457, 463], $\text{Bi}_2\text{Sr}_{1.6}\text{Gd}_{0.4}\text{CuO}_{6+\delta}$ [338], $(\text{Cu,C})\text{Ba}_2\text{Ca}_3\text{Cu}_4\text{O}_{12+\delta}$ [440], and $\text{TlBa}_2\text{Ca}_2\text{Cu}_2\text{O}_{10-\delta}$ [442]. Nanoscale electronic nonhomogeneity on the crystal surface was shown to substantially affect the CDW-like DOS modulation observed by STM in $\text{Bi}_2\text{Sr}_{1.4}\text{La}_{0.6}\text{CuO}_{6+\delta}$ [340].

Large gap scatterings obviously do not correlate with sharp transitions into the superconducting state at any doping of well prepared samples (implying Cooper-pairing

homogeneity), which was demonstrated, for example, by specific heat studies [464]. To solve the problem, one should bear in mind that the gaps measured by STM technique are of two kinds (in our opinion, superconducting gaps and pseudogaps—CDW gaps), which cannot be easily distinguished experimentally [81–83, 160, 337]. The guess was proved in [291], where contributions of both gaps in the STM spectra of $(\text{Bi}_{0.62}\text{Pb}_{0.38})_2\text{Sr}_2\text{CuO}_{6+x}$ were separated by an ingenious trick. Namely, the authors normalized the measured local conductances by removing the larger-gap inhomogeneous background. Then, it became clear that the superconducting gap is more or less homogeneous over the sample's surface, whereas the larger gap (the pseudogap, i.e., the CDW gap) is essentially inhomogeneous.

The intimate origin of the pseudogap variations is currently not understood. At the same time, the inhomogeneity of electron characteristics is also inherent to the related solid solutions BPB, which was demonstrated by spatially resolved electron energy loss spectroscopy [465]. It is reasonable to suggest that this inhomogeneity both in BPB and high- T_c oxides is strengthened near free surfaces in agreement with Josephson current measurements across BPB bicrystal tunnel boundaries [466].

Still, there is an interesting phenomenon, which might explain trends for electric properties in cuprates to be inhomogeneous. We mean a spontaneous phase separation, suggested long ago for antiferromagnets [467–470] and the electron gas in paramagnets [471–474]. This idea was later transformed into stripe activity in cuprate and manganite physics, where alternating conducting and magnetic regions constituted separated “phases” [12, 302, 475–480]. Recently, a lot of evidence for local lattice distortions, Jahn-Teller polaron occurrence, and other percolation and filamentary structure formation appeared [177, 217, 228, 481–485], supporting new sophisticated theoretical efforts in the science of phase separation [84, 230, 379, 486–493], mostly but not necessarily dealing with high- T_c oxides. The electronic inhomogeneity in cuprates, as discussed above, belongs to the same category of phenomena. Whatever its origin, intrinsic inhomogeneity of cuprates and other oxides seems to be an important feature that needs explanation in order to understand superconductivity (much more homogeneous) itself. Note that electronic phase separation into magnetic and nonmagnetic domains was also found in the iron pnictide superconductor $\text{Ba}_{1-x}\text{K}_x\text{Fe}_2\text{As}_2$ [494], whereas disorder-induced inhomogeneities of superconducting properties was observed in TiN films [495].

Another high- T_c oxide, $\text{YBa}_2\text{Cu}_3\text{O}_{6+x}$, containing CuO chains in addition to CuO_2 planes, was known for a long time as a material exhibiting one-dimensional CDWs [496]. However, the authors of more recent tunnel measurements [497] concluded that the would-be CDW manifestations might have a different nature, since the observed one-dimensional modulation wavelengths have rather a strong dispersion. Nevertheless, it seems that in view of the large CDW amplitude scatter in BSCCO discovered later, this conclusion is premature, with local variations of the FS shape being a possible origin of CDW wave vector modifications.

As one sees from the evidence discussed above, CDW modulations are observed in cuprates both directly (as patterns of localized energy-independent electron states in the conventional $\mathbf{r}$ -space) and indirectly (as concomitant gapping phenomena). The pseudogap energy $E_{\text{PG}} > \Delta_{\text{SC}}$ constitutes an appropriate scale for CDW gapping. Here, Δ_{SC} is the superconducting gap. On the other hand, at low energies $E < \Delta_{\text{SC}}$, single-particle tunneling spectroscopy probes mixed electron-hole d -wave Bogoliubov quasiparticles [498], which are delocalized excitations. In this case, it is natural to describe the tunnel conductance in the momentum, $\mathbf{k}$ -space. The interference between Bogoliubov quasiparticles is especially strong for certain wave vectors $\mathbf{q}_i$ ($i = 1, \dots, 16$) connecting extreme points on the constant energy contours [499–502]. The interference $\mathbf{k}$ -space patterns involve those wave vectors [343, 416, 499, 503–505], this picture being distinct from and complementary to the partially disordered CDW unidirectional or checkerboard structures [344, 359, 365, 371, 458, 506–508].

It is remarkable that interference $\mathbf{r}$ -space patterns on cuprate surfaces, the latter being in the superconducting state, are not detected, contrary to the clear-cut STM observations of electron de Broglie standing waves, induced by point defects or step edges, revealed in conductance maps on the normal metal surfaces [509, 510]. The latter waves are in effect Friedel oscillations [511] formed by two-dimensional normal electron density crests and troughs with the wave length π/k_F , k_F being the Fermi wave vector. On the other hand, spatial oscillating structures of local DOS in the d -wave superconducting state are determined by other representative vectors $\mathbf{q}_i$, so that the characteristic oscillations can be denominated as Friedel-like ones at most [502, 512]. Nevertheless, the attenuation of both kinds of spatial oscillations due to superconducting modifications of the screening medium should be more or less similar. Namely, in the isotropic superconducting state, the electron gas polarization operator loses its original singularity at $k = 2k_F$ for gapped FS sections [513]. As a consequence, Friedel oscillations gain an extra factor $\exp(-2r/\pi\xi_0)$ [514, 515], where ξ_0 is the BCS coherence length [498]. For d -wave superconductors, the attenuation will be weaker and will totally disappear in the order-parameter node directions. However, those distinctions are not crucial, since the nodes have a zero measure. The modification of screening by formation of Bogoliubov quasiparticles in d -wave high- T_c oxides explains the absence of conspicuous spatial structures in STM maps, which correspond to the wave vectors $\mathbf{q}_i$ mentioned above.

We consider the observed CDWs in oxides as a consequence of electron-hole (dielectric) pairing on the nested sections of corresponding FSs [158–160, 516]. Such a viewpoint is also clearly supported by the experiments in layered dichalcogenides [374–376], the materials analogous to cuprates in the sense of superconductivity appearance against the dielectric (CDW) partial gapping background [378, 379]. At the same time, other sources of CDW instabilities are also possible [517, 518]. As for the microscopic mechanism causing CDW formation, it might be an electron-phonon (Peierls insulator) [519, 520] or a Coulomb

one (excitonic insulator) [431, 521, 522], or their specific combination. Excitonic instability may also lead to the SDW state [522, 523], also competing with superconductivity for the FS [160, 524–529]. It should be noted that researchers asserted that they found plenty of Peierls insulators or partially gapped Peierls metals [158–160, 530–532]. At the same time, the excitonic phase, being mathematically identical in the mean-field limit [533] and physically similar [534] to the Peierls insulator, was not identified unequivocally. One can only mention that some materials claimed to be excitonic insulators, namely, a layered transition-metal dichalcogenide 1T-TiSe₂ with a commensurate CDW [535, 536], alloys TmSe_{0.45}Te_{0.55} [537], Sm_{1-x}La_xS [538], and Ta₂NiSe₅ with a *direct band gap* at the Brillouin zone Γ point in the parent high- T state [539]. Therefore it is reasonable that precisely in the later case, the low- T excitonic state is not accompanied by CDWs.

It is necessary to indicate that in many cases, the claimed “charge stripe order” and the more unpretentious “charge order” are an euphemism describing the old good CDWs: “Stripes is a term that is used to describe unidirectional density-wave states, which can involve unidirectional charge modulations (charge stripes) or coexisting charge and spin-density order spin stripes” [12]. We do not think it makes sense to use the term “stripes” in the cases of pure CDW or spin-density-wave (SDW) ordered states. At the same time, this term should be reserved for different possible more general kinds of microseparation [12, 477, 479, 540–542], having nothing or little to do with periodic lattice distortions, FS nesting, or Van Hove singularities. The need to avoid misnomers and duplications while naming concepts is quite general in science, as was explicitly stressed by John Archibald Wheeler, who himself coined many terms in physics (“black hole” included) [543].

In this connection, it seems that some experimentalists unnecessarily vaguely attribute the spatially periodical charge structure in the low-temperature tetragonal phase of La_{1.875}Ba_{0.125}CuO₄, revealed by X-ray scattering [544], to the hypothetical nematic structure or the checkerboard Wigner crystal. Indeed, quite similar spatial charge structures found in La_{1.875}Ba_{0.125-x}Sr_xCuO₄ by neutron scattering [545] were correctly and without reservation identified as CDW-related ones, whereas a checkerboard structure (if any) can be considered as a superposition of two mutually perpendicular CDWs. The same can be written about the “stripe” terminology used in [546], where X-ray scattering revealed a periodical charge structure in the low-temperature tetragonal phase of another cuprate La_{1.8-x}Eu_{0.2}Sr_xCuO₄.

One should mention two other possible collective states competing with Cooper pairing. Namely, these are states with microscopic orbital and spin currents that circulate in the ground state of excitonic insulator (there can be four types of the latter [522]). The concept of the state with current circulation, preserving initial crystal lattice translational symmetry, was invoked to explain cuprate properties [547]. Another order parameter, hidden from clear-cut identification by its supposed extreme sensitivity to sample imperfection, is the so-called d -density wave-order parameter [548, 549]. It is nothing but a CDW order parameter times the same

form-factor $f(k) = \cos(k_x) - \cos(k_y)$, the product being similar to that for $d_{x^2-y^2}$ -superconductors. Here, k_x and k_y are the wave-vector components in the CuO_2 plane. To some extent, the dielectric order parameter of the Bilbro-McMillan model [159, 160, 385] and its generalizations—they are presented below—contains the same physical idea as in the d -density-wave model: nonuniformity of the CDW gap function in the momentum space.

Although the destructive CDW action on superconductivity of many good materials is beyond question [56, 160, 380, 384, 545, 550, 551], it *does not mean* that maximal T_c are limited by *this factor only*. For instance, T_c falls rapidly with the hole concentration p in overdoping regions of $T_c - p$ phase diagrams for different Pb-substituted $\text{Bi}_2\text{Sr}_2\text{CuO}_{6+\delta}$ compounds, even in the case when the critical doping value p_{cr} corresponding to $T_d \rightarrow 0$ lies outside the superconducting dome [552]. A Cu-doped superconducting chalcogenide Cu_xTiSe_2 constitutes another example confirming the same trend [380]. Namely, CDW manifestations die out for $x \gtrsim 0.06$, whereas T_c starts to decrease for $x > x_{\text{optimal}} = 0.08$. As has been already mentioned, overdoping can reduce T_c simply owing to screening of matrix elements for electron-phonon interaction [99, 215, 216].

4. Theory of CDW Superconductors and Its Application to Cuprates

The majority of our results presented below were obtained for s -wave superconductors with CDWs. It is a case, directly applicable to many materials (e.g., dichalcogenides, trichalcogenides, tungsten brozes, etc.). On the other hand, as was indicated above, the exact symmetry of the superconducting order parameter in cuprates is not known, although the d -wave variant is considered by most researchers in the field as the ultimate truth. Notwithstanding any future solution of the problem, our theory of CDW-related peculiarities in quasiparticle tunnel CVCs can be applied to cuprates, since we are not interested in small energies $eV < \Delta$, where the behavior of a reconstructed DOS substantially depends on whether it is the s - or d -wave order parameter [553–555]. Here, $e > 0$ is the elementary charge, and Δ is the amplitude of the superconducting order parameter. As for the thermodynamics of CDW superconductors, we present both s - and d -cases, each of them having their own specific features.

4.1. Thermodynamics of s -Wave CDW Superconductors. The Dyson-Gorkov equations for the normal ($\mathcal{G}_{ij}$) and anomalous ($\mathcal{F}_{ij}$) temperature Green's functions in the case of coupled superconducting $\Delta_{ij}^{\alpha\gamma}$ and dielectric (CDW) $\Sigma_{ij}^{\alpha\gamma}$ matrix order parameters are the starting point of calculations and can be found elsewhere [160, 386, 397, 426, 427]. Greek superscripts correspond to electron spin projections, and italic subscripts describe the natural split of the FS into degenerate (nested, d) and non-degenerate (non-nested, n) sections. For the quasiparticles on the nested sections, the standard condition leading to the CDW gapping holds:

$$\xi_1(\mathbf{p}) = -\xi_2(\mathbf{p} + \mathbf{Q}), \quad (1)$$

where $\mathbf{p}$ is the quasimomentum, $\mathbf{Q}$ is the CDW vector (see the discussion above), Planck's constant $\hbar = 1$. This equation binds the electron and hole bands $\xi_{1,2}(\mathbf{p})$ for the excitonic insulator [431, 522] and different parts of the one-dimensional self-congruent band in the Peierls insulator case [516]. At the same time, the rest of the FS remains undistorted below T_d and is described by the electron spectrum branch $\xi_3(\mathbf{p})$. Such an approach was suggested long ago by Bilbro and McMillan [385]. We adopt the strong-mixing approximation for states from different FS sections. This means an appearance of a single superconducting order parameter for d and nd FS sections. The spin-singlet structure (s -wave superconductivity and CDWs) of the matrix normal ($\Sigma_{ij}^{\alpha\gamma} = \Sigma\delta_{\alpha\beta}$) and anomalous ($\Delta_{ij}^{\alpha\gamma} = I_{\alpha\beta}$) self-energy parts (where $(I_{\alpha\beta})^2 = -\delta_{\alpha\beta}$) in the weak-coupling limit is suggested. Here, $\delta_{\alpha\beta}$ is the Kronecker delta. The self-consistency equations for the order parameters obtained in accordance with the fundamentals can be expressed in the following form [386]:

$$\begin{aligned} 1 &= V_{\text{BCS}}N(0)[\mu I(D) + (1 - \mu)I(\Delta)], \\ 1 &= V_{\text{CDW}}N(0)\mu I(D), \end{aligned} \quad (2)$$

where

$$I(x) = \int_0^\Omega \frac{d\xi}{\sqrt{\xi^2 + x^2}} \tanh \frac{\sqrt{\xi^2 + x^2}}{2T}. \quad (3)$$

Here, the Boltzmann constant $k_B = 1$, V_{BCS} and V_{CDW} are contact four-fermion interactions responsible for superconductivity and CDW gapping, respectively. The gap

$$D(T) = [\Delta^2(T) + \Sigma^2(T)]^{1/2} \quad (4)$$

is a combined gap appearing on the nested FS sections, whereas the order parameter Δ defines the resulting observed gap on the rest of the FS (compared with the situation in cuprates [344, 350, 356, 403]). The parameter μ characterizes the degree of the FS dielectrization (hereafter, we use this nonconventional term instead of “gapping” in some places to avoid confusion with the superconducting gapping), so that $N_d(0) = \mu N(0)$ and $N_{nd}(0) = (1 - \mu)N(0)$ are the electron DOSs per spin on the FS for the nested and nonnested sections, respectively. The upper limit in (3) is the relevant cut-off frequency, which is assumed to be equal for both interactions. If the cut-offs BCS and CDW are considered different, the arising correction, $\log(\Omega_{\text{CDW}}/\Omega_{\text{BCS}})$, is logarithmically small [385] and does not change qualitatively the subsequent results. Only in the case of almost complete electron spectrum dielectric gapping ($\mu \rightarrow 1$) does the difference between BCS and CDW become important for the phase coexistence problem [433]. This situation is, however, of no relevance for substances with detectable superconductivity, since T_c tends to zero for $\mu \rightarrow 1$. In this subsection, we confine ourselves to the case $\text{Re}\Sigma > 0$, $\text{Im}\Sigma = 0$, since the phase φ of the complex order parameter $\Sigma \equiv |\Sigma|e^{i\varphi}$ does not affect the thermodynamic properties, whereas tunnel currents *do depend* on φ [160, 556, 557], which will be demonstrated explicitly below.

Introducing the bare order parameters $\Delta_0 = 2\Omega \exp[-1/V_{\text{BCS}}N(0)]$ and $\Sigma_0 = 2\Omega \exp[-1/V_{\text{CDW}}N_d(0)]$, we can rewrite the system of (2) in an equivalent form, convenient for numerical calculations:

$$\begin{aligned} I_M[\Delta, T, \Delta(0)] &= 0, \\ I_M(D, T, \Sigma_0) &= 0, \end{aligned} \quad (5)$$

where

$$I_M(G, T, G_0) = \int_0^\infty \left(\frac{1}{\sqrt{\xi^2 + G^2}} \tanh \frac{\sqrt{\xi^2 + G^2}}{2T} - \frac{1}{\sqrt{\xi^2 + G_0^2}} \right) d\xi \quad (6)$$

is the standard Mühlischlegel integral [558], the root of which $G = s\text{Mü}(G_0, T)$ is the well-known gap dependence for the s -wave BCS superconductor [9], $G_0 = G(T=0)$, and [385]

$$\Delta(0) = (\Delta_0 \Sigma_0^{-\mu})^{1/(1-\mu)}. \quad (7)$$

However, (5) mean that both gaps $\Delta(T)$ and $D(T)$ have the BCS form $G = s\text{Mü}(G_0, T)$ [386], namely: (i) $\Delta(T) = s\text{Mü}[\Delta(0), T]$, that is, the actual value of the superconducting gap of the CDW superconductor at $T=0$ is $\Delta(0)$ rather than Δ_0 , and the actual superconducting critical temperature is $T_c = \gamma\Delta(0)/\pi$; (ii) at the same time, $D(T) = s\text{Mü}(\Sigma_0, T)$, which determines $T_d = \gamma\Sigma_0/\pi$. Here, $\gamma = 1.7810\dots$ is the Euler constant.

From (4), we obtain that, at $T=0$,

$$\Sigma_0^2 = \Delta^2(0) + \Sigma^2(0). \quad (8)$$

Replacing $\Delta(0)$ by its value (7), we arrive at the conclusion that in the model of s -wave superconductor with partial CDW gapping, two order parameters coexist only if $\Delta_0 < \Sigma_0$. Then, according to (7), $\Delta(0) < \Delta_0$; that is, the formation of the CDW, if it happens, always inhibits superconductivity, in agreement with the totality of experiments [160, 375, 380, 382, 551]. Also, vice versa, according to (4), for $T < T_c$, $\Sigma(T) < s\text{Mü}(\Sigma_0, T)$; that is, superconductivity suppresses dielectrization.

In Figure 3, the dependences $\Delta(T)$ and $\Sigma(T)$ are shown for various parameters of the partially dielectrized CDW s -wave superconductor. It can be easily inferred from the data shown in both panels that, in agreement with the foregoing, $\Delta(T)/\Delta(0)$ curves coincide with the Mühlischlegel one for any values of the dimensionless parameters μ and $\sigma_0 \equiv \Sigma_0/\Delta_0$. The novel feature, which has been overlooked in other investigations, is the possibility of such a strong suppression of Σ for low enough T that it becomes *smaller* than Δ , although T_d is larger than T_c (see Figure 3(b)). This intriguing situation can be realized for the parameter σ_0 close to unity. One should note that the actual gaps Δ and D (the former coincides with the superconducting order parameter) are monotonic functions of T . However the dielectric order parameter is not.

The magnitudes of T_c and $\Delta(0)$ strongly depend on μ and σ_0 , although the simple BCS-like scaling between them

survives, that is, for CDW s -wave superconductors $\Delta(0)/T_c = \pi/\gamma \approx 1.76$. Although for, say, $\Sigma_0 \geq 1.5\Delta_0$ and reasonable $\mu = 0.5$ [386], the demand of self-consistency between $\Sigma(T)$ and $\Delta(T)$ becomes less important quantitatively. It justifies our previous approach with T -independent Σ [427] and the estimation of combined gap as $(\Delta_{\text{BCS}}^2(T) + \Delta_{\text{PG}}^2)^{1/2}$ with T -independent Δ_{PG} made on the basis of interlayer tunneling measurements in BSCCO mesas [559]; self-consistency leads to new qualitative effects and cannot be avoided. As for the magnitude of the very Δ_{PG} , inferred from tunneling measurements, it was found in [559] to be substantially smaller than that of $\Delta_{\text{BCS}}(T \rightarrow 0)$, whereas the opposite case turned out to be true both for BSCCO [349, 399, 560, 561], $\text{Bi}_{2-x}\text{Pb}_x\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ [460], and $(\text{Bi,Pb})_2\text{Sr}_2\text{Ca}_2\text{Cu}_3\text{O}_{10+\delta}$ [562]. Other tunnel measurement for BSCCO [417] revealed $\Delta_{\text{PG}} > \Delta_{\text{BCS}}(T \rightarrow 0)$ for underdoped samples and $\Delta_{\text{PG}} < \Delta_{\text{BCS}}(T \rightarrow 0)$ for overdoped ones. A marked sensitivity of Δ_{PG} to doping together with strong inhomogeneity, discovered in Bi-based ceramics [334–336, 338, 343, 359, 440, 441, 456–458, 563, 564] and $\text{Ca}_{2-x}\text{Na}_x\text{CuO}_2\text{Cl}_2$ [565], may be responsible for the indicated discrepancies.

Since the BCS character of the gap dependences for the CDW s -wave superconductor is preserved, the T -dependence of the heat capacity C for the doubly gapped electron liquid (i.e., below the actual T_c) equals to the superposition of two BCS-like functions:

$$C(T) = \frac{2\pi^2 N(0)}{3} \left[(1-\mu) T_c c_{\text{BCS}}\left(\frac{T}{T_c}\right) + \mu T_d c_{\text{BCS}}\left(\frac{T}{T_d}\right) \right], \quad (9)$$

where

$$c_{\text{BCS}}\left(t = \frac{T}{T_c^{\text{BCS}}}\right) = \frac{C_{\text{BCS}}(T)}{C_{\text{BCS}}(T = T_c^{\text{BCS}} + 0)}. \quad (10)$$

It should be noted that the normalized discontinuity $\Delta C/C_n(T_c)$ at the superconducting phase transition may also serve as indirect evidence for the CDW gap on the FS, because in this case it is *not at all* a trivial BCS jump:

$$\frac{\Delta C_{\text{BCS}}}{\gamma_S T_c} = \frac{12}{7\zeta(3)} \approx 1.43. \quad (11)$$

Here, $C_n(T) = \gamma_S T \equiv (2\pi^2 N(0)/3)T$ is the normal electron-gas heat capacity, whereas γ_S is the Sommerfeld constant. CDW-driven deviations from the BCS behavior was recognized long ago [425, 566]. However, only the self-consistent approach [386] allows us to give a quantitative answer at any value of the parameters appropriate to the partially CDW-gapped superconductor. It can be seen from Figure 4(a), where the conventionally normalized superconducting phase transition anomaly is shown as a function of μ . The discontinuity is always smaller than the BCS value (11), in agreement with previous qualitative considerations [425, 566]. At the same time, the BCS ratio is restored not only for $\mu = 0$, that is, in the absence of the dielectrization, but also for $\mu \rightarrow 1$. In the former case, it is clear, because we are dealing with a conventional BCS superconductor. On the other hand, for large enough μ , CDW gapping

FIGURE 3: Temperature dependences of the superconducting (Δ) and dielectric (Σ) order parameters for different values of the dimensionless parameters μ (the portion of the nested Fermi surface sections, where the charge-density-wave, CDW, gap develops) and σ_0 (see explanations in the text). (Taken from [386].)

almost completely destroys superconductivity, so $T_c \ll T_d$. Therefore, in the relevant superconducting T range, the contribution to $C(T)$ from the d FS sections, governed by the gap $D \approx \Sigma$, becomes exponentially small. Another term, determined by the n FS section, ensures the BCS limiting value of the normalized discontinuity.

The dependences of $\Delta C/C_n$ on σ_0 for various values of μ are depicted in Figure 4(b). One sees that the effect is large for σ_0 close to unity, whereas the difference between 1.43 and $\Delta C/C_n$ goes to zero as σ_0^{-2} , verifying the asymptotical result [425]. It should be noted that the heat capacity calculation scheme adopted for s -wave CDW superconductors can be applied also to other types of order parameter symmetry.

Experimental data on heat capacity, which could confirm the expressed ideas, are scarce. For Nb_3Sn , it was recently shown by specific heat measurements using the thermal relaxation technique that $T_c \approx 17 \div 18$ K is reduced when the critical temperature of the martensitic transition $T_d \approx 42 \div 53$ K grows [567]. Unfortunately, a large difference between T_c and T_d made the effects predicted by us quite small here, which is probably the reason why they have not been observed in these studies.

As for cuprates, reference should be made to $\text{La}_{2-x}\text{Ba}_x\text{CuO}_{4-y}$ [568], $\text{La}_{2-x}\text{Sr}_x\text{CuO}_{4-y}$ [569], and $\text{YBa}_2\text{Cu}_3\text{O}_{7-y}$ [570], where underdoping led to a reduction of $\Delta C/C_n$. The same is true for measurements of the heat capacity in $\text{Bi}_2\text{Sr}_{2-x}\text{La}_x\text{CuO}_{6+\delta}$ single crystals [571], which demonstrated that the ratio $\Delta C/C_n$ for a strongly underdoped sample turned out to be about 0.25, that is, much below BCS values $12/7\zeta(3) \approx 1.43$ and $8/7\zeta(3) \approx 0.95$ [572] for s -wave and d -wave superconductivity, respectively.

There is also an opposite evidence for the relationship $\Delta C/C_n > 1.43$, for example, in the electron-doped high- T_c oxide $\text{Pr}_{1.85}\text{Ce}_{0.15}\text{CuO}_{4-\delta}$ [573]. More details, as well as information on other CDW superconductors, can be found in [386]. In any case, despite the well-known challenging controversy for BPB [392, 574–576], the problem was not studied enough for any superconducting oxide family, probably due to experimental difficulties.

4.2. Enhancement of the Paramagnetic Limit in s -Wave CDW Superconductors. Upper critical magnetic fields H_{c2} [577–579] (along with critical currents [132, 579, 580]) belong to main characteristics of superconductors crucial for their applications. In particular, knowing the upper limits on upper critical fields is necessary to produce superconducting materials for high-performance magnets, not to talk about scientific curiosity.

One of such limiting factors is the paramagnetic destruction of spin-singlet superconductivity, which was discovered long ago theoretically by Clogston [581] and Chandrasekhar [582]. In the framework of the BCS theory, they obtained a limit

$$H_p^{\text{BCS}} = \frac{\Delta_{\text{BCS}}(T=0)}{\mu_B^* \sqrt{2}} \quad (12)$$

from above on H_{c2} at zero temperature, T . Here, $\Delta_{\text{BCS}}(T)$ is the energy gap in the quasiparticle spectrum of BCS s -wave superconductor, and μ_B^* is the effective Bohr magneton, which may not coincide with its bare value $\mu_B = e\hbar/2mc$, where $\hbar$ is Planck's constant, equal to unity in the whole

FIGURE 4: Dependences of the normalized heat capacity discontinuity ΔC at T_c on μ (a) and σ_0 (b). (Taken from [386].)

article but shown here explicitly for clarity, m is the electron mass, and c is the velocity of light.

Limit (12) may be overcome at a high concentration of strong spin-orbit scatterers, when the spins of the electrons, constituting the Cooper pairs, are flipped [583]. Then, the actual $H_{c2}(T = 0)$ starts to exceed [584] the classical bound. Such an enhancement of H_{c2} has been observed, for example, in Al films coated with Pt monolayers [585]. The Pt atoms served there as strong spin-orbit scatterers due to their large nuclear charge Z . One should indicate a possibility of exceeding value (12), if the energy, E , dependence of the normal state density of states is significant, which is the case in the neighborhood of the van Hove singularity [518]. Then, the BCS approximation of $N(E) \simeq N(0)$ is no longer valid, so that the actual H_p may become larger than limit (12) [586].

We have found another reason, why the Clogston-Chandrasekhar value can be exceeded. Namely, it is the appearance of a partial CDW-driven dielectric gap on the d sections of the FS [427, 587–589]. The expected increase of the calculated limiting paramagnetic field H_p for CDW superconductors, as compared to H_p^{BCS} , is intimately associated with paramagnetic properties of the normal CDW state, which are very similar to those for BCS s -wave superconductors [382, 590–592].

It should be emphasized that the very self-consistency of the two-gap solution [386] made the treatment of the paramagnetic limit problem [589] transparent and less involved than previous approximations [427, 587, 588].

To calculate the paramagnetic limit, we considered the relevant free energies F per unit volume for all possible ground state phases in an external magnetic field H . The parent nonreconstructed phase (actually existing only above T_d !), with both superconducting and CDW pairings switched

off and in the absence of H , served as a reference point. At $T < T_d$, we deal with relatively small differences δF reckoned from this hypothetical “doubly-normal” state [498]. In our case, in the Clogston-Chandrasekhar spirit [581, 582], there are two energy differences to be compared [589], specifically, that of a paramagnetic phase in the magnetic field [593] (diamagnetic effects are not taken into account when one is interested in the paramagnetic limit *per se*)

$$\delta F_p = -N(0)(\mu_B^* H)^2 \quad (13)$$

and that of a CDW-superconducting phase

$$\delta F_s = -N_n(0) \frac{\Delta^2(0)}{2} - N_d(0) \frac{D^2(0)}{2}. \quad (14)$$

Here, $\Delta(0)$ is determined by (7), whereas $D(0)$, as stems from (8), is equal to $\Sigma_0 = \pi T_d/\gamma$. A simple algebra leads to the analytical equation for the increase of the paramagnetic limit over the Clogston-Chandrasekhar value (12):

$$\left(\frac{H_p}{H_p^{\text{BCS}}} \right)^2 = 1 + \mu \left[\left(\frac{\Sigma_0}{\Delta_0} \right)^{2/(1-\mu)} - 1 \right]. \quad (15)$$

This relationship is expressed in terms of genuine (bare) system parameters μ , Σ_0 , and Δ_0 . However, experimentalists are interested in the dependence of H_p/H_p^{BCS} on actual measurable quantities. The transformation of (15) can be easily made, and one arrives at the final formula

$$\begin{aligned} \left(\frac{H_p}{H_p^{\text{BCS}}} \right)^2 &= 1 + \mu \left[\left(\frac{\Sigma(0)}{\Delta(0)} \right)^2 - 1 \right] \\ &= 1 + \mu \left[\left(\frac{T_d}{T_c} \right)^2 - 1 \right]. \end{aligned} \quad (16)$$

FIGURE 5: Contour plot of the ratio H_p/H_p^{BCS} on the plane $(T_c/T_d, \mu)$. Here H_p is the paramagnetic limit for CDW superconductors and H_p^{BCS} is that for BCS spin-singlet superconductors, T_c and T_d are the observed critical temperatures of the superconducting and CDW transitions, respectively. (Taken from [589].)

To calculate the expected paramagnetic limit, one needs to know μ , which was estimated, for example, as 0.2 for NbSe_3 [594] or 0.15 for $\text{La}_{2-x}\text{Sr}_x\text{CuO}_4$ [595].

The contour curves in the parameter space obtained from (16) are displayed in Figure 5. One can readily see that for typical $T_c/T_d \approx 0.05\text{--}0.2$ (some A15 compounds are rare exceptions [160]) and moderate values of $\mu \approx 0.3\text{--}0.5$, the augmentation of the paramagnetic limit becomes very large. There is a number of CDW superconductors [589], where the increase of the paramagnetic limit was detected, in accordance with the results presented here. Unfortunately, not so much can be said about high- T_c oxides. It seems that extremely high values of H_{c2} observed in these materials are the reason of the unjustified neglect to the problem.

4.3. Dip-Hump Structures and Pseudogaps in Tunnel Current-Voltage Characteristics for Junctions Involving CDW Superconductors. In Section 3, a lot of evidence was presented concerning dip-hump structures (DHSs) and pseudogaps in high- T_c oxides [81–83, 160, 559]. Broadly speaking, pseudogaps and DHSs have much in common. In particular, they can coexist with superconducting coherent peaks, their appearance in current-voltage characteristics (CVCs) is to some extent random, and their shapes are sample-dependent. Therefore, a suggestion inevitably arises that those two phenomena might be governed by the same mechanism. Our main assumption is that both pseudogaps and dip-hump structures are driven by CDW instabilities discussed above and that their varying appearances are coupled with the intrinsic, randomly inhomogeneous electronic structure of cuprates. In the strict sense, according to the adopted scenario, both DHSs and pseudogaps are the manifestations of the *same* dielectric DOS depletion,

the former being a result of superimposed CDW- and superconductivity-induced CVC features below T_c . To justify our approach, it is crucial that direct spatial correlations between irregular patterns of CDW three-dimensional supermodulations [365] and topographic maps of the superconducting gap amplitudes on the BSCCO surface were displayed by tunneling spectroscopy [458].

A detailed description of our approaches to the problems of tunneling through junctions with CDW isotropic superconductors as electrodes, and the emergence of DHSs in the CVCs of high- T_c oxides can be found elsewhere [81–83, 160, 386, 596]. Here, we will present only a summary of our new results, briefly touching only those issues that are necessary for the rest of the paper to be self-contained.

We should emphasize different roles of the order parameter phases in determining quasiparticle tunnel currents. Concerning the superconducting order parameter, its phase may be arbitrary unless we are interested in the Josephson current across the junction. On the other hand, the CDW phase φ governs quasiparticle CVCs of junctions with CDW superconductors as electrodes [556, 557]. The value of φ can be pinned by various mechanisms in both excitonic and Peierls insulators, so that φ acquires the values either 0 or π in the first case [522] or is arbitrary in the unpinned state of the Peierls insulator [516]. At the same time, in the case of an inhomogeneous CDW superconductor, which will be discussed below, a situation can be realized, where φ values are not correlated over the junction area. Then, the contributions of elementary tunnel currents may compensate one another to some extent, and this configuration can be phenomenologically described by introducing a certain effective phase φ_{eff} of the CDW order parameter. If the spread of the phase φ is random, the most probable value for φ_{eff} is $\pi/2$, and the CVC for a nonsymmetric junction involving CDW superconductor becomes symmetric.

Most often, CVCs for cuprate-I-N (i.e., S-I-N) junctions reveal a DHS only at $V = V_S - V_N < 0$ [597–599], so that the occupied electron states below the Fermi level are probed for CDW superconductors. In our approach, it corresponds to the phase φ close to π . This preference may be associated with some unidentified features of the CDW behavior near the sample surface.

On the other hand, there are S-I-N junctions, where DHS structures are similar for both V polarities [347, 442, 443, 600]. As for those pseudogap features, which were unequivocally observed mostly at high T , no preferable V -sign of their manifestations was found. We note that the symmetry of the tunnel conductance $G(V) = dJ/dV$, where J is the tunnel current through the junction, might be due either to the microscopic advantage of the CDW state with $\varphi = \pi/2$ or to the superposition of different current paths in every measurement covering a spot with a linear size of a CDW coherence length at least. Both possibilities should be kept in mind. The variety of $G(V)$ patterns in the S-I-N set-up for the same material and with identical doping is very remarkable, showing that the tunnel current is rather sensitive to the CDW phase φ . Nevertheless, the very appearance of the superconducting domain structure for cuprates with local domain-dependent gaps and critical

temperatures [601] seems quite plausible for materials with small coherence lengths. Essentially the same approach has been proposed earlier to explain superconducting properties of magnesium diboride [78].

On the basis of information presented above and using the self-consistent solutions for $\Delta(T)$ and $\Sigma(T)$, we managed to describe the observed rich variety of $G(V)$ patterns by calculating quasiparticle tunnel CVCs $J(V)$ for two typical experimental set-ups. Namely, we considered S-I-N and S-I-S junctions, where “S” here means a CDW superconductor. A unique tunnel resistance in the normal state R enters into all equations, since we assume the incoherent tunneling to occur, in accordance with the previous analysis for BSCCO [122, 602]. The used Green’s function method followed the classical approach of Larkin and Ovchinnikov [603]. We skip all (quite interesting) technical details, since they can be found elsewhere [81–83, 160, 386, 596].

The obtained equations for $J(V)$ form the basis for calculations both $J(V)$ and $G(V)$ (sub-, superscripts ns and s denote S-I-N and S-I-S junctions, resp.). They must be supplemented with a proper account of the nonhomogeneous background, since, as was several times stressed above, STM maps of the cuprate crystal surfaces consist of random nano-scale patches with different gap depths and widths, as well as coherent edge sharpnesses. In this connection, our theory assumes the combination CDW + inhomogeneity to be responsible for the appearance of the DHSs. Our main conclusion is that it is the dispersion of the parameter Σ_0 —and, as a result, the D -peak smearing (the Δ -peak also becomes smeared but to a much lesser extent)—that is the most important to reproduce experimental pictures. The value of the FS gapping degree μ is mainly responsible for the amplitude of the DHSs. At the same time, neither the scattering of the parameter μ nor that of the superconducting order parameter Δ_0 can result in the emergence of smooth DHSs, so that sharp CDW features remain unaltered. Therefore, for our purpose, it was sufficient to average only over Σ_0 rather than simultaneously over all parameters of CDW superconductors, although the variation of any individual parameter made the resulting theoretical CVCs more similar to experimental ones.

Although it is a well-recognized matter of fact that CDW-driven D -singularities in $G(V)$ scatter more strongly for a nonhomogeneous medium than main coherent superconducting peaks at $eV = \pm\Delta$ (S-I-N junctions) or $\pm 2\Delta$ (S-I-S junctions), this phenomenon has not yet been explained. It seems that the sensitivity of the Peierls [604, 605] or excitonic-insulator [431, 606] order parameters to the Coulomb potential of the impurities, for example, oxygen ions, might be the reason of such a dispersion. On the other hand, s -wave superconductivity is robust against impurity influence (Anderson theorem [607–609]). As for anisotropic superconductivity with d -wave or other kinds of symmetry, they are suppressed by nonmagnetic impurity scattering [3, 554, 610, 611] due to scattering-induced order parameter isotropization. Their survival in disordered cuprate samples, especially in the context of the severe damage inflicted by impurities on the pseudogap, testifies that the Cooper-pair order parameter includes a substantial isotropic component.

The parameter Σ_0 was assumed to be distributed within the interval $[\Sigma_0 - \delta\Sigma_0, \Sigma_0 + \delta\Sigma_0]$. The normalized weight function $W(x)$ was considered as a bell-shaped fourth-order polynomial within this interval and equal to zero beyond it (see the discussions in [81]). In any case, the specific form of $W(x)$ is not crucial for the final results and conclusions.

Our approach is in essence the BCS-like one. It means, in particular, that we do not take a possible quasiparticle “dressing” by impurity scattering and the electron-boson interaction, as well as the feedback influence of the superconducting gapping, into account [612, 613]. Those effects, important *per se*, cannot qualitatively change the random two-gap character of superconductivity in cuprates.

As was already mentioned, we have assumed so far that both Δ and Σ are s -wave-order parameters. Nevertheless, our approach to CVC calculations is qualitatively applicable to superconductors with the d -wave symmetry, if not to consider the intragap voltage range $|eV| < \Delta$.

The results of calculations presented below show that the same CDW + inhomogeneity combination can explain DHSs at low T as well as the pseudogap phenomena at high T , when the DHS is smoothed out. Thus, theoretical T -dependences of tunnel CVCs mimic the details of the DHS transformation into the pseudogap DOS depletion for nonsymmetric and symmetric junctions, involving cuprate electrodes. We consider the CDW-driven phenomena, DHS included, as the tip of an iceberg, a huge underwater part of which is hidden by strong superconducting manifestations, less influenced by randomness than their CDW counterpart. To uncover this part, one should raise T , which is usually done with no reference to the DHS, the latter being substantially smeared by the Fermi-distribution thermal factor. It is this DOS depletion phenomenon that is connected to the pseudogapping phenomena [14, 18, 334, 348, 399, 441].

The results of calculations of $G^{ns}(V)$ in the case when parameter Σ_0 is assumed to scatter are shown in Figure 6 for $\varphi = \pi$. The value $\varphi = \pi$ was selected, because this case corresponds to the availability of the DHS in the negative-voltage branch of the nonsymmetric CVC, and such an arrangement is observed in the majority of experimental data. In accordance with our basic equations, all the four existing CVC peculiarities at $eV = \pm\Delta$ and $\pm D$ become smeared, although to various extent: the large singularities at $eV = \pm\Delta$ almost preserve their shape, the large singularity at $eV = -D$ transforms into a DHS, and the small one at $eV = D$ disappears on the scale selected. The one-polarity dip-hump peculiarity in experimental CVCs for BSCCO [597] is reproduced excellently. Owing to relationship (7), the actual parameter Δ also disperses, but, due to the small value of μ , this fluctuation becomes too small to be observed in the plot. Thus, the calculated CVCs of Figure 6 demonstrate all principal features intrinsic to the tunnel conductivity of S-I-N junctions at low T , involving CDW superconductors, specifically, asymmetry with respect to the V sign is associated with the phase $\varphi \neq \pi/2$ of the CDW-order parameter, the emerging CDW induces singularities at $eV = \pm D$, whereas the intrinsic CDW inhomogeneity transforms the major one into a DHS, totally suppressing the minor.

FIGURE 6: (Color online) Bias voltage, V , dependences of the dimensionless differential conductance $RG(V) = R dJ/dV$ for the tunnel junction between an inhomogeneous CDW superconductor and a normal metal, expressed in energy units. Here, J is the quasiparticle tunnel current, R is the resistance of the junction in the normal state, and $e > 0$ is the elementary charge. The bare parameters of the CDW superconductor are $\Delta_0^* = 20$ meV, $\Sigma_0^* = 50$ meV, and $\mu = 0.1$; the temperature $T = 4.2$ K. Various dispersions $\delta\Sigma_0^*$ centered around the mean value $\Sigma_0^* = 50$ meV. (Taken from [81].)

An example of the transformation, with T , of the DHS-decorated tunnel spectra into the typical pseudogap-like ones is shown in Figure 7 for S-I-N junctions with $\varphi = \pi$ (panel (a)) and $\pi/2$ (panel (b)). The CDW-superconductor parameters are $\Delta_0 = 20$ meV, $\Sigma_0 = 50$ meV, $\mu = 0.1$, and $\delta\Sigma_0 = 20$ meV; the temperature $T = 4.2$ K. For this parameter set, the “actual” superconducting critical temperatures T_c of random domains lie within the interval of 114–126 K, and T_d is in the range of 197–461 K. From Figure 7(a), the transformation of the DHS-including pattern of the CVCs calculated for $T \ll T_c$ into the pseudogap-like ones in the vicinity of T_c or above it becomes clear. The asymmetric curves displayed in (a) are similar to the measured STM $G^{ns}(V)$ dependences for overdoped and underdoped BSCCO compositions [441]. The overall asymmetric slope of the experimental curves, which is independent of gaps and T , constitutes the main distinction between them and our theoretical results. It might be connected to the surface charge carrier depletion induced by CDWs and mentioned above. Another interesting feature of our results is a modification and a shift of the Δ -peak. Although Δ diminishes as T grows, the Δ -peak moves toward higher bias voltages; such a behavior of the Δ -peak is to be undoubtedly associated with its closeness to the Σ -governed DHS. In experiments, a confusion of identifying this Δ -driven singularity with a pseudogap feature may arise,

since the observed transformation of Δ -features into D -ones looks very smooth [348].

It is notable that in the case of asymmetric $G^{ns}(V)$, the low- T asymmetry preserves well into the normal state, although the DHS as such totally disappears. The extent of the sample randomness substantially governs CVC patterns. Therefore, pseudogap features might be less or more pronounced for the same materials and doping levels. At the same time, for the reasonable spread of the problem parameters, the superconducting coherent peaks always survive the averaging (below T_c , of course), in accordance with experiment. Our results also demonstrate that the dependences $\Delta(T)$ taken from the tunnel data may be somewhat distorted in comparison to the true ones due to the unavoidable Δ versus Σ interplay. One should stress that in our model, “hump” positions, which are determined mainly by Σ rather than by Δ , anticorrelate with true superconducting gap values Δ inferred from the coherent peaks of $G(V)$. It is exactly what was found for nonhomogeneous BSCCO samples [614].

Similar CDW-related features should be observed in the CVCs measured for symmetric S-I-S junctions. The $G^s(V)$ dependences for this case with the same sets of parameters as in Figure 7 are shown in Figure 8. In analogy with symmetric junctions between BCS superconductors, one would expect an appearance of singularities at $eV = \pm 2\Delta$, $\pm(D + \Delta)$, and $\pm 2D$. Such, indeed, is the case. However, the magnitudes of the features are quite different (the details of the analysis can be found in [81, 83]). As readily seen, the transformation of the symmetric DHS pattern into the pseudogap-like picture is similar to that for the nonsymmetric junction. This simplicity is caused by a smallness of the parameter $\mu = 0.1$, so that the features at $eV = \pm 2D$, which are proportional to μ^2 , are inconspicuous on the chosen scale. At the same time, the singularities at $eV = \pm(D + \Delta)$ are of the *square-root type*. Note that for *arbitrary* Σ - and Δ -magnitudes, those energies do not coincide with the values $\pm(\Sigma + \Delta)$ (in more frequently used notation, $\pm(\Delta_{PG} + \Delta_{SG})$), which can be sometimes met in literature [615]. The later relation becomes valid only for $\Sigma \gg \Delta$.

The appearance of the T -driven zero-bias peaks is a salient feature of certain CVCs displayed in Figure 8. As is well known [603], this peak is caused by tunneling of thermally excited quasiparticles between empty states with an enhanced DOS located above and below equal superconducting gaps in symmetric S-I-S junctions. Such a feature was found, for example, in $G^s(V)$ measured for grain-boundary symmetric tunnel junctions in epitaxial films of the s -wave oxide CDWS $\text{Ba}_{1-x}\text{K}_x\text{BiO}_3$ [616]. One should be careful not to confuse this peak with the dc Josephson peak restricted to $V = 0$, which is often seen for symmetric high- T_c junctions [399]. The distinction consists in the growth of the quasiparticle zero-bias maximum with increasing T up to a certain temperature, followed by its drastic reduction. On the other hand, the Josephson peak decreases monotonously as $T \rightarrow T_c$.

The profile and the behavior of the zero-bias peak at nonzero T can be explained in our case by the fact that, in effect, owing to the nonhomogeneity of electrodes,

FIGURE 7: $G(V)$ dependences for the tunnel junction between an inhomogeneous CDWS and a normal metal for various temperatures T . The CDW order parameter phase $\varphi = \pi$ (a) and $\pi/2$ (b), and the spread of the CDW order parameter-amplitude $\delta\Sigma_0^* = 20$ meV. All other parameters are indicated in the text. (c) STM spectra for underdoped BSCCO-Ir junctions registered at various temperatures. (Reprinted from [598], taken from [83].)

the junction is a combination of a large number of *symmetric* and *nonsymmetric* junctions with varying gap parameters. The former compose a mutual contribution to the current in the vicinity of the $V = 0$ point, and the width of this contribution along the V -axis is governed by temperature alone. On the other hand, every junction from the latter group gives rise to an elementary current peak in the CVC at a voltage equal to the relevant gap difference. All such elementary contributions form something like a hump around the zero-bias point, and the width of this hump along the V -axis is governed by the sum of actual—dependent on

the zero- T values and on the temperature itself—gap spreads in both electrodes. It is clear that the T -behavior of the current contribution of either group is rather complicated, to say nothing of their combination.

From our CVCs calculated for both nonsymmetric (Figures 6 and 7) and symmetric (Figure 8) junctions, it comes about that the “dip” is simply a depression between the hump, which is mainly of the CDW origin, and the superconducting coherent peak. Therefore, as has been noted in [617], the dip has no separate physical meaning. It disappears as T increases, because the coherent peak forming

FIGURE 8: (a) The same as in Figure 7(a), but for a symmetric junction between similar CDW superconductors. (b) Temperature variations of experimental differential current-voltage characteristics (CVCs) for a $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ break junction. (Reprinted from [399], taken from [83].)

the other shoulder of the dip fades down, so that the former dip, by expanding to the $V = 0$ point, becomes an integral constituent of the shallow pseudogap minimum.

Therefore, it became clear that the CDW manifestations against the nonhomogeneous background can explain both *subtle* DHS structures in the tunnel spectra for high- T_c oxides and *large* pseudogap features observed both below and above T_c . The DHS is gradually transformed into the pseudogap-like DOS, lowering as T grows. Hence, both phenomena are closely interrelated, being in essence the manifestations of the same CDW-governed feature smeared by inhomogeneity of CDW superconductors. Therefore, the DHS and pseudogap features should not be treated separately. The dependences of the calculated CVCs on the CDW phase φ fairly well describe the variety of asymmetry manifestations in the measured tunnel spectra for BSCCO and related compounds.

4.4. Coexistence of CDWs and d -Wave Superconductivity. We recognize that some of our results, which were obtained assuming that the superconducting order parameter coexisting with CDWs is isotropic, might be applicable to cuprates with certain reservations, since a large body of evidence in favor of $d_{x^2-y^2}$ symmetry in high- T_c oxides [131, 132, 305, 306, 618, 619] is available, although there are experimentally-based objections [109, 116–129]. In any case, it seems

instructive to extend the partial dielectrization approach to d -wave Cooper pairing. For simplicity, we argue in terms of two-dimensional first Brillouin zone and Fermi surface, neglecting c -axis quasiparticle dispersion, which should be taken into account, in principle [620]. Since the dielectric, Σ , and superconducting, d -wave Δ , order parameters have different momentum dependences, their joint presence in the electron spectrum is no longer reduced to a combined gap (4), as it was for isotropic superconductivity.

In the d -wave case, superconductivity is described by a weak-coupling model with a Hamiltonian given, for example, in [553, 621]. In accordance with photoemission [371, 622–624] and STM [359, 368, 370, 506, 507, 512, 625] data (see Figure 1), the mean-field CDW Hamiltonian is restricted to momenta near flat-band regions, antinodal from the viewpoint of the four-lobe d -wave gap-function $\Delta(T) \cos 2\theta$ [306]. In those regions, the nesting conditions (1) between pairs of mutually coupled quasiparticle branches are fulfilled. For instance, static CDW wave vectors $\mathbf{Q} = (\pm 2\pi/4.2a_0, 0)$ and $(0, \pm 2\pi/4.2a_0)$ —with an accuracy of 15%—in $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ are revealed in STM studies [359]. Thus, we characterize a CDW checkerboard state (symmetric with respect to $\pi/2$ -rotations) by four sectors in the momentum space centered with the lobes and with an opening 2α each ($\alpha < \pi/4$). It should be noted that vectors $\mathbf{Q}$ depend on doping, which was explicitly shown

FIGURE 9: (Color online) Order-parameter maps for a conventional d -wave superconductor (Δ , solid curve) and a partially gapped CDW metal (Σ , dashed curve).

for $\text{Bi}_2\text{Sr}_2\text{CuO}_{6+\delta}$ [343]. The dielectric (CDW-induced) order parameter is $\Sigma(T)$ inside the 2α -cones, being angle-independent here, and zero outside (see Figure 9).

The plausibility of this scenario is supported—at least partially—by recent STM studies of intrinsically inhomogeneous BSCCO samples [626]. Specifically, the authors analyzed the composition, temperature, and angular dependences of the gaps on various FS sections and showed that nodal superconducting gaps for overdoped specimens exhibit more or less conventional d -wave behavior, whereas in underdoped samples nodal (superconducting) and antinodal gaps (CDW gaps, as is assumed here) superimpose on one another in tunnel spectra. It is important that for underdoped compositions antinodal gaps do not change drastically with T , when crossing T_c . The conclusion made in [626] that the entire FS contributes to bulk superconductivity in overdoped samples corresponds—if proved to be correct—to the actual shrinkage of nested FS sections, that is, to $\mu \rightarrow 0$.

We obtained a new set of Dyson-Gor'kov equations for normal and superconducting Green's functions for the system with electron-hole of whatever nature and d -wave Cooper pairings, which were solved in the same straightforward manner as in the s -wave case [160, 386] (see above). We arrived at the system of two coupled equations for $\Delta(T)$ and $\Sigma(T)$:

$$\int_0^{\mu\pi/4} I_M(\sqrt{\Sigma^2 + \Delta^2 \cos^2 2\theta}, T, \Sigma_0) d\theta = 0, \quad (17)$$

$$\begin{aligned} \int_0^{\mu\pi/4} I_M(\sqrt{\Sigma^2 + \Delta^2 \cos^2 2\theta}, T, \Delta_0) \cos^2 2\theta d\theta \\ + \int_{\mu\pi/4}^{\pi/4} I_M(\Delta \cos 2\theta, T, \Delta_0) \cos^2 2\theta d\theta = 0, \end{aligned} \quad (18)$$

where $\mu = 4\alpha/\pi$ is the dielectrically gapped portion of the FS for the specific model of partial gapping, shown in Figure 9, and $I_M(\Delta \cos 2\theta, T, \Delta_0)$ is the Mühschlegel integral (6). The analysis of the generic T - δ phase diagram for cuprates shows that both Σ_0 and μ reduce with doping, whereas the hole-like FS pockets centered at the $(\pi/a_0, \pi/a_0)$ point of the Brillouin zone shrink for every specific high- T_c oxide (see, e.g., [343]).

On the other hand, in the absence of CDW gapping, (18) becomes a d -wave gap equation:

$$\int_0^{\pi/4} I_M(\Delta \cos 2\theta, T, \Delta_0) \cos^2 2\theta d\theta = 0, \quad (19)$$

the solution of which $\Delta = d\text{Mü}(\Delta_0, T)$ is known [553, 621]. In particular, the critical temperature is $T_{c0} = (2\Omega\gamma/\pi) \exp[-1/V_{\text{BCS}}N(0)]$, as in the s -wave case. From (19), it follows that in agreement with [553], $(\Delta_0/T_{c0})_d = (2/\sqrt{e})(\pi/\gamma)$, revealing a modified “ d -wave” BCS-ratio different from the s -pairing value

$$\left(\frac{\Delta_0}{T_{c0}}\right)_s = \frac{\pi}{\gamma} \approx 0.824 \left(\frac{\Delta_0}{T_{c0}}\right)_d. \quad (20)$$

Here, e is the base of natural logarithm. It is evident that our model takes into account many-body correlations both explicitly (the emergence of two pairings) and implicitly (via the renormalization of the parameters Σ_0 and μ). Weak-coupling values of the ratio Δ_0/T_{c0} for other anisotropic order parameter symmetries do not differ much from the value of $(\Delta_0/T_{c0})_d$ [627, 628].

Due to the different order parameter symmetry, readily seen from (17) and (18), the situation is mathematically more involved than for isotropic CDW superconductors, where a simple relationship (4) takes place. This was not recognized in a recent work [629], where the opposite wrong statement was made. *Prima facie* subtle mathematical differences between descriptions of s -wave and d -wave CDW superconductors lead to conspicuous physical consequences. Indeed, the numerical dependences $\Delta(T)$ and $\Sigma(T)$ found from (17) and (18) and shown in Figure 10 differ *qualitatively* from their counterparts $\Delta_s(T)$ and $\Sigma_s(T)$ in a certain range of model parameters. (In this subsection, we do not introduce a natural subscript “ d ” for brevity.) Figure 10 (a) demonstrates that a reduction of the bare parameter Σ_0 , keeping Δ_0 and μ constant, resulting in the transformation of $\Sigma(T)$ with a cusp at $T = T_c$ and a concave region at $T < T_c$ (the behavior appropriate for CDW s -superconductors in the whole allowable parameter range, as is demonstrated in Figure 3) into curves describing a *novel* peculiar reentrant CDW state. It is remarkable that the reentrance found by us is appropriate to an extremely simple basic model with two competing order parameters. At the same time, the CDW structures in real systems may be much more complicated with nonmonotonic T -dependencies even in the absence of superconductivity [352].

Let us formulate conditions necessary to observe this crossover. First, (20) means that $\Delta(T)/\Delta_0$ for conventional d -superconductors is steeper than $(\Delta(T)/\Delta_0)_s$. In our case, it means that $\Delta(T)/\Delta_0$, when the CDW disappears, is steeper than $\Sigma(T)/\Sigma_0$ in the absence of superconductivity, which is described by (6). Hence, for the CDW phase to exist (the upper critical temperature $T_{\text{CDW}}^u > 0$), it should be $T_{\text{CDW}}^u = (\gamma/\pi)\Sigma_0 > T_{c0} = (\sqrt{e}\gamma/2\pi)\Delta_0$. As a consequence, the first constraint on the model parameters should be fulfilled: $\Sigma_0 > (\sqrt{e}/2)\Delta_0 \approx 0.824\Delta_0$. The constraint stems from the competition between emerging Δ and Σ on the d FS section only. The actual coexistence between superconductivity and

FIGURE 10: (Color online) Temperature, T , dependences of the normalized (a) CDW Σ and (b) superconducting Δ gap functions. Δ_0 equal to $\Delta(T = 0)$ when CDWs are absent is 1. The values of Σ_0/Δ_0 equal to $\Sigma(T = 0)/\Delta_0$ in the absence of superconductivity are 1.5 (1), 1.2 (2), 1 (3), 0.95 (4), 0.9 (5), 0.85 (6), and 0.8 (7); $\mu = 0.3$.

CDWs was not involved in these reasonings, so the inequality does not include the control parameter μ . Therefore, T_{CDW}^u thus defined coincides with $T_{\text{CDW}0}$.

Second, below the lower critical temperature of the CDW reentrance region, T_{CDW}^l , if any, (18) defines $\Delta(T) = d\text{M}\ddot{\text{u}}(\Delta_0, T)$, and we should use (17) with $T = T_{\text{CDW}}^l$ and $\Delta(T_{\text{CDW}}^l) = d\text{M}\ddot{\text{u}}(\Delta_0, T_{\text{CDW}}^l)$ to determine $T_{\text{CDW}}^l(\Delta_0, \Sigma_0, \mu)$ numerically. The crossover value of Σ_0^{cr} , when $T_{\text{CDW}}^l = 0$, corresponds to the separatrix on the order parameter- T plane, dividing possible $\Sigma(T)$ -curves (see Figure 10(a)) into two types: reentrant and nonreentrant. However, (17) brings about $\Sigma_0^{\text{cr}} = \Delta_0 \exp[(4/\mu\pi) \int_0^{\mu\pi/4} \ln(\cos 2\theta) d\theta]$. To observe the reentrant behavior, the second constraint should be $\Sigma < \Sigma_0^{\text{cr}}$. For the curves in Figure 10, $\mu = 0.3$ was chosen, so that we obtain the reentrance range $0.824\Delta_0 < \Sigma_0 < 0.963\Delta_0$, which agrees with numerical solutions of the full self-consistent equation set. We emphasize that CDWs survives the competition with d -wave superconductivity even at $\Sigma_0/\Delta_0 < 1$, which is not the case for stronger isotropic Cooper pairing (see the discussions above).

In Figure 10(b), the concomitant $\Delta(T)$ dependences are depicted. One sees how d -wave superconductivity, suppressed at large Σ_0 , recovers in the reentrance parameter region. Therefore, two regimes of CDW manifestation can be observed in superconductors. In both cases, the CDW is seen as a pseudogap above T_c [81, 83] in photoemission and tunnel experiments. However, the corresponding DHS at low T may either be observed or not, depending on whether the reentrance occurs. This might be an additional test for an anisotropic (not necessarily $d_{x^2-y^2}$ -wave) Cooper pairing to dominate in cuprates.

To control the change-over between different regimes in cuprates, one can use either hydrostatic pressure or doping. In both cases, μ is the main varying parameter. In Figure 11,

the curves $\Sigma(T)$ and $\Delta(T)$ are shown for $\Sigma_0/\Delta_0 = 0.9$ and various μ . It is readily seen how drastic is the low- T depression of Σ by superconductivity, when the dielectrically gapped FS sectors are small enough. Doping $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ [403] and $(\text{Bi,Pb})_2(\text{Sr,L a})_2\text{CuO}_{6+\delta}$ [356] with oxygen was shown to sharply shrink the parameter μ . Note that the $\Delta(T)$ dependences are distorted by CDWs, and they do *not* coincide with the scaled “parent” curve— $d\text{M}\ddot{\text{u}}(T)$, in this case—in contrast to what should be observed for CDW s -superconductors (Figure 3). Therefore, various observed forms of $\Delta(T)$ *per se* cannot unambiguously testify to the superconducting pairing symmetry. Moreover, cuprate superconductivity might be, for example, a mixture of s - and d -wave contributions [130, 630].

It is evident that different strengths of CDW-imposed suppression of the superconducting energy gap in the electron spectrum Δ and the critical temperature T_c must change the ratio $\Delta(0)/T_c$ —the benchmark of *weak-coupling* superconductivity (see (20)). If one recalls that, as was shown above, this ratio in CDW s -superconductors remains the same as in conventional s -ones, the situation becomes very intriguing. In Figure 12(a), the dependences of $2\Delta(0)/T_c$ and T_c/Δ_0 ratios on Σ_0/Δ_0 are displayed. One sees that $2\Delta(0)/T_c$ sharply increases with Σ_0/Δ_0 for $\Sigma_0/\Delta_0 \leq 1$ and swiftly saturates for larger Σ_0/Δ_0 , whereas T_c/Δ_0 decreases almost evenly. The saturation value proves to be 5.2 for $\mu = 0.3$. We stress that such large enhancement of $2\Delta(0)/T_c$ agrees well with experimental data [441, 478, 631, 632] for cuprates and *cannot* be achieved taking into account strong-coupling electron-boson interaction effects for reasonable relationships between T_c and effective boson frequencies ω_E [633, 634] (one can hardly accept, e.g., the value $T_c/\omega_E \approx 0.3$ [634] as practically meaningful). Furthermore, the destruction of the alternating-sign superconducting order

FIGURE 11: (Color online) The same as in Figure 10 but for $\Sigma_0/\Delta_0 = 0.9$ and $\mu = 0.1$ (1), 0.3 (2), 0.5 (3), 0.6 (4).

parameter by impurity scattering approximated by collective boson modes also could not explain [635] high values of $2\Delta(0)/T_c$, for example, inherent to underdoped BSCCO [347, 597]. Therefore, our weak-coupling model is *sufficient* to explain—*on its own*—the large magnitude of $2\Delta(0)/T_c$ in cuprates, possible strong-coupling effects resulting in at most minor corrections.

Another possible alternative reason of high $2\Delta(0)/T_c$ ratios might be a singular energy dependence of the normal-state electron DOS near the FS, for instance, near the Van Hove anomalies in low-dimensional electron subsystems [518]. It turned out, however, that, at least in the weak-coupling (BCS) approximation for *s*-wave Cooper pairing, the ratio $2\Delta(0)/T_c$ is not noticeably altered [636, 637]. Moreover, calculations in the framework of the strong-coupling Eliashberg theory [10] showed that the van Hove singularity influence on T_c is even smaller than in the BCS limit [638]. Furthermore, weak-coupling calculations for orthorhombically distorted hole-doped cuprate superconductors (without CDWs) demonstrated that $2\Delta(0)/T_c$ can be estimated as an intermediate between *s*-wave and *d*-wave limits [639], being smaller than needed to explain the experiment. It means that our approach remains so far the only one capable of explaining high $2\Delta(0)/T_c \approx 5 \div 8$ (and even larger values [632]) for cuprates. We emphasize that it is very important to reconcile theoretical values for $2\Delta(0)/T_c$ as well as $\Delta C/\gamma_S T_c$ with experimental ones. Otherwise, the microscopic theory becomes “too” phenomenological with Δ/T_c as an *additional free parameter* of the system [640].

It is instructive from the methodological point of view to mention a previous unsuccessful attempt to explain the increase of $2\Delta(0)/T_c$ by a pseudogap influence [641]. The authors of this reference assumed the *identical d*-wave symmetry for both the superconducting, $\Delta(T)$, and

temperature-independent pseudogap, E_{PG} , order parameters. Additionally, dielectric gapping was supposed to be effectively complete rather than partial, the latter being intrinsic to our model and follows from the experiments for cuprates. These circumstances excluded self-consistency imposed on E_{PG} , namely, $E_{PG} \lesssim 0.53\Delta_0(T=0)$, where $\Delta_0(T=0)$ is the parent superconducting order parameter amplitude. At the same time, it is well known that for existing CDW superconductors the strength of CDW instability is at least not weaker than that of its Cooper-pairing counterpart [160]. We should emphasize once more that the *main peculiarity* of our model, dictated by the observations, which led to the adequate description of thermodynamic properties for *d*-wave superconductors with CDWs, is the distinction between relevant order parameter symmetries.

The μ -dependences of $2\Delta(0)/T_c$ and T_c/Δ_0 are shown in Figure 12(b). They illustrate that $2\Delta(0)/T_c$ can reach rather large values, if the dielectric gapping sector is wide enough. This growth is however limited by a drastic drop of T_c leading to a quick disappearance of superconductivity. We think that it is exactly the case of underdoped cuprates, when a decrease of T_c is accompanied by a conspicuous widening of the superconducting gap. For instance, such a scenario was clearly observed in break-junction experiments for $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ samples with a large doping range [642].

As was pointed out in [478], various photoemission and tunneling measurements for different cuprate families show a typical average value $2\Delta(0)/T_c \approx 5.5$. From Figure 12(b), we see that this ratio corresponds to $\mu \approx 0.35$ at $\Sigma_0/\Delta_0 = 1$. The other curve readily gives $T_c/\Delta_0 \approx 0.35$. Since $\Delta_0/T_{c0} \approx 2.14$ for a *d*-wave superconductor (see above), we obtain $T_c/T_{c0} \approx 0.75$, being quite a reasonable estimation of T_c -reduction by CDWs.

FIGURE 12: (Color online) Dependences of $2\Delta(0)/T_c$ (squares) and T_c/Δ_0 (circles) on Σ_0/Δ_0 (panel (a), $\mu = 0.3$) and μ (panel (b), $\Delta_0/\Sigma_0 = 1$). T_c is the superconducting critical temperature, d -BCS ≈ 4.28 is a value for a conventional superconductor with d -wave symmetry of the order parameter.

4.5. SDWs and Superconductivity. There are plenty of materials, where SDWs compete with superconductivity, although a simultaneous existence of the order parameters cannot be always proved [160, 643, 644]. For completeness, we give here certain short comments on the latest developments in this direction.

During the last years, an interest arose to the phase with the hidden order parameter in URu_2Si_2 , emerging at about 17.5 K and being some kind of SDWs, coexisting with superconductivity at $T < 1.5$ K [153, 645–647]. It should be noted that the partial gapping idea applied to SDW materials [397, 422–424, 566, 648–650] was invoked to explain ordering in this compound at the times of the discovery [651].

More attention was paid to Cr and its alloys, where CDWs and SDWs are linked and coexist [652, 653]. It might be interesting to observe mutual influence of CDWs and SDWs on superconductivity [654–657].

The problem of an interplay between SDWs and superconductivity received strong impetus recently, especially because of the fundamental discovery of magnetic element-based high- T_c pnictide superconductors [49, 654–657]. Theoretical efforts were also continued (see, e.g., [658–662]). It is worthwhile noting that, for certain doping ranges, superconducting cuprates also demonstrate [663] the coexistence of Cooper pairing with SDWs rather than CDWs, the latter being appropriate for the majority of high- T_c oxide compositions (see above and [160]).

Although the coexistence of superconductivity with SDWs or more exotic orbital antiferromagnetic and spin current ordering [522, 548, 664] is left beyond the scope of this review, the relevant physics is not less fascinating than that of their CDW-involving analogues.

5. Conclusions

The presented material testifies that CDWs play the important role in high- T_c oxides and govern some of the properties that usually have been considered as solely determined by superconductivity *per se*. Sometimes CDWs manifest themselves explicitly (observed checkerboard or unidirectional structures, DHSs, pseudogaps) but, in the majority of phenomena, they “only”—but often drastically—change the magnitude of certain effects in the superconducting state (the heat capacity anomaly, the paramagnetic limit, the T -dependence of H_{c2} , the $\Delta(0)/T_c$ ratio). Cuprates are not unique as materials with coexisting CDWs and superconductivity, but the scale of the interplay is very large here due to the strength of the Cooper pairing in those compounds.

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